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THE VELOCITY OF SOUND IN LIQUID HELIUM NEAR THE λ POINT*

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Recent measurements of the velocity and attenuation of sound in liquid helium[1] showed that just below the λ point there exists a relaxation time τ proportional to $(T_\lambda - T)^{-1}$, in agreement with a theoretical conclusion of Landau and Khalatnikov.[2] Although this fact is of intrinsic interest, the presence of relaxation with its attendant velocity dispersion obscures the behavior of the equally interesting quantity u_0 (the velocity of sound in the low-frequency limit $\omega\tau \ll 1$). Indeed, since τ becomes infinite at the transition temperature, measurements of u_0 very close to T_λ could be made only at an inconveniently low frequency. A method of calculating the temperature dependence of u_0 near the λ point is described in this Letter.

The calculation is based on the idea, suggested by Pippard,[3] that near a second-order phase transition the entropy and volume surfaces, considered as functions of p and T, are cylindrical sections whose generators are parallel to the transition line. This "cylindrical approximation" leads to the following relations among the second derivatives of the Gibbs function:

$$C_p = \alpha V T_\lambda \beta + C_0, \tag{1}$$

and

$$\beta = \alpha \kappa_T + \beta_0, \tag{2}$$

where C_p is the specific heat at constant pressure, β the coefficient of thermal expansion, κ_T the isothermal compressibility, and α the slope of the λ line, $(dp/dT)_\lambda$. C_0 and β_0 are given by

$$C_0 = \alpha T_\lambda (dS/dp)_\lambda, \quad \beta_0 = (\alpha/V)(dV/dp)_\lambda. \tag{3}$$

Equation (1) has been found experimentally to hold over a range of about ±0.040°K about the λ point,[4] and it is reasonable to suppose that Eq. (2) has approximately the same range of validity. It is evident from the cylindrical form of the entropy and volume surfaces that derivatives taken along any line parallel to the λ line in the p, T plane (that is, along a generator of the surface) will have the same value irrespective of the point on the surface at which they are evaluated. The meaning of the subscript λ can therefore be generalized to include such parallel lines. Derivatives such as those appearing in Eqs. (3) then become formally functions of p and T, but as long as the cylindrical approximation holds may be regarded as constants.

The adiabatic compressibility may now be expressed in terms of the above quantities. Choosing p and S as independent variables,

$$\left(\frac{\partial V}{\partial p}\right)_S = \left(\frac{dV}{dp}\right)_\lambda - \left(\frac{\partial V}{\partial S}\right)_p \left(\frac{dS}{dp}\right)_\lambda, \tag{4}$$

and

$$\left(\frac{\partial T}{\partial p}\right)_S = \left(\frac{dT}{dp}\right)_\lambda - \left(\frac{\partial T}{\partial S}\right)_p \left(\frac{dS}{dp}\right)_\lambda. \tag{5}$$

By means of the Maxwell relation $(\partial V/\partial S)_p = (\partial T/\partial p)_S$ and Eq. (5), the quantity $(\partial V/\partial S)_p$ can be eliminated from Eq. (4), yielding

$$\left(\frac{\partial V}{\partial p}\right)_S = \left(\frac{dV}{dp}\right)_\lambda - \left(\frac{dT}{dp}\right)_\lambda \left(\frac{dS}{dp}\right)_\lambda + \left(\frac{\partial T}{\partial S}\right)_p \left(\frac{dS}{dp}\right)_\lambda^2. \tag{6}$$

Using Eqs. (3) and neglecting the small difference between T and T_λ, this may be written in the

form

$$\left(\frac{\partial V}{\partial p}\right)_S = \frac{V\beta_0}{\alpha} + \frac{C_0}{T_\lambda \alpha^2}\left(\frac{C_0}{C_p} - 1\right). \quad (7)$$

Finally, the square of the sound velocity is

$$u_0^2 = -V^2\left(\frac{\partial V}{\partial p}\right)_S^{-1} = -\frac{T_\lambda V^2 \alpha^2}{\alpha V T_\lambda \beta_0 + C_0(C_0/C_p - 1)}. \quad (8)$$

C_0 can be evaluated from experimental data on the coefficient of thermal expansion[4] and specific heat[5] by means of Eq. (1), and β_0 can be chosen to fit the observed sound velocity at a temperature sufficiently far below T_λ for dispersion to be unimportant. Equation (8) then determines the temperature dependence of u_0 near the λ point. This is illustrated in Fig. 1, where the full curve has been calculated from Eq. (8) and the broken curve reproduces the experimental results at 1 Mc/sec.[1] The scale of the ordinate refers to the calculated curve, which has been fitted to the experimental one 2 millidegrees below the λ point. The agreement is adequate below T_λ; the rather large discrepancy above the λ point may indicate that the cylindrical approximation is not valid in this region, or may be the result of experimental errors in the more difficult helium I range.

An important property of the cylindrical approximation is that it should become more exact as the transition temperature is approached. Accordingly, the qualitative behavior in the immediate vicinity of the λ point is of particular significance. There are three essential features of this behavior: (i) u_0 is continuous at the λ point, (ii) it falls to a finite value at the transition temperature, and (iii) its rate of change becomes infinite as the λ point is approached from either side. This is a direct consequence of the behavior of C_p, which is the only temperature-dependent factor in Eq. (8).

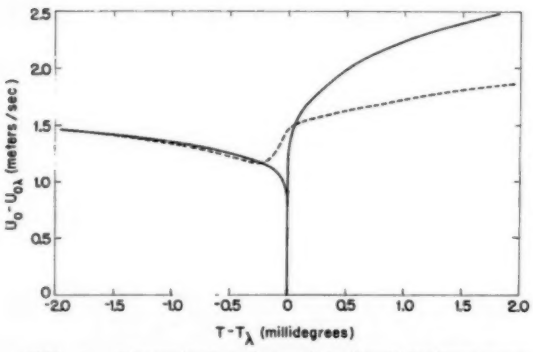

FIG. 1. The velocity of sound in liquid helium near the λ point. —Calculated [Eq. (8)]; --- observed at 1 Mc/sec (reference 1).

It is instructive to compare Eq. (8) with the expression for the velocity discontinuity at the λ point which can be derived from the Ehrenfest relations[2]:

$$\Delta\left(\frac{1}{u_0^2}\right) = -\frac{T_\lambda}{V^2}\Delta\left(\frac{1}{C_p}\right)\left(\frac{dS}{dp}\right)_\lambda^2. \quad (9)$$

[This equation can also be deduced directly from Eq. (8).] Equation (9) clearly shows that the velocity discontinuity must vanish if the specific heat becomes infinite at the transition temperature, in agreement with (i), but gives no information about the temperature dependence of u_0 as the λ point is approached. Such information can only be obtained by means of some specific assumption about the shape of the entropy and volume surfaces, such as Pippard's cylindrical approximation.

It is perhaps worth mentioning that the absence of a velocity discontinuity raises the following difficulty in Landau and Khalatnikov's calculation of the attenuation of sound near the λ point.[2] The magnitude of the sound absorption is determined in the usual manner by the velocity dispersion immediately below T_λ, that is, by the difference between the velocity at infinite frequency, u_∞, and that at zero frequency, u_0. According to Landau and Khalatnikov, u_0 is simply the equilibrium velocity of sound in helium II, while u_∞ is equal to the velocity in helium I. The quantity $u_\infty - u_0$ is therefore equal to the velocity discontinuity at the λ point. If this discontinuity vanishes there can be no dispersion, and consequently no excess sound absorption. Since large excess absorption does indeed occur, it is clear that Landau and Khalatnikov's approach is inadequate in the presence of a specific heat singularity.

In conclusion, it should be pointed out that Eq. (8) is not restricted to liquid helium, and should be equally valid for any phase transition near which Eqs. (1) and (2) hold.

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PROPERTIES OF A SEMICONDUCTOR SURFACE AS DETERMINED FROM A MODIFIED DRIFT MOBILITY EXPERIMENT

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In connection with a study of the role of semiconductor surface layers in the behavior of metal to semiconductor contacts, an attempt has been made to observe directly any injection-extraction effects in the semiconductor bulk due only to the semiconductor surface. Montgomery[1] has already assumed that such effects occur in explaining some of his field-effect results observed at high frequencies.

In the present experiments, the drift mobility technique[2] is employed with a modification in the emitter circuit as shown in Fig. 1. Capacitive coupling is made to the semiconductor surface by means of a 1 mm × 2 mm brass electrode separated from the surface by 10^{-2} cm of strontium titanate.[3] The capacity of this condenser is about 10^{-11} farad. Current will flow to or from the surface only when the emitter pulse is turned on or off. During these surges, carrier disturbances may be produced in the semiconductor. The nature of these disturbances can be examined and analyzed if they are made to drift past a collector electrode, as is generally done in the drift mobility experiment.

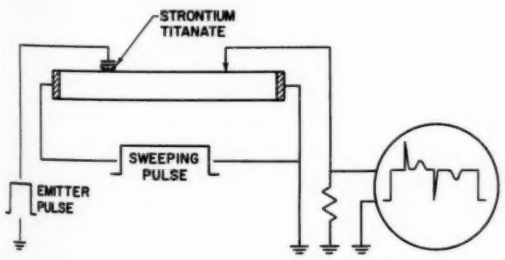

FIG. 1. Modified drift mobility technique for studying semiconductor surfaces.

Signals have been observed in this experiment for etched, ground, and diamond polished surfaces on germanium samples. Photographs of some typical signals are shown in Fig. 2 for a 35 ohm-cm, p-type sample. These signals are interpreted as actually resulting from the injection and extraction effects of the surface.

It is helpful to compare the injection-extraction effects of the surface-bulk combination to those of normal n-p or p^+-p type junctions.[4] For example, if the surface is n-type with respect to the bulk, injection will occur in the bulk when a negative pulse is turned on, thereby inducing a positive surface charge; extraction will occur when the pulse is turned off. The detailed mechanism of how this might happen is outlined by Montgomery.[1] It is also helpful to refer to the surface conductivity versus surface barrier po-

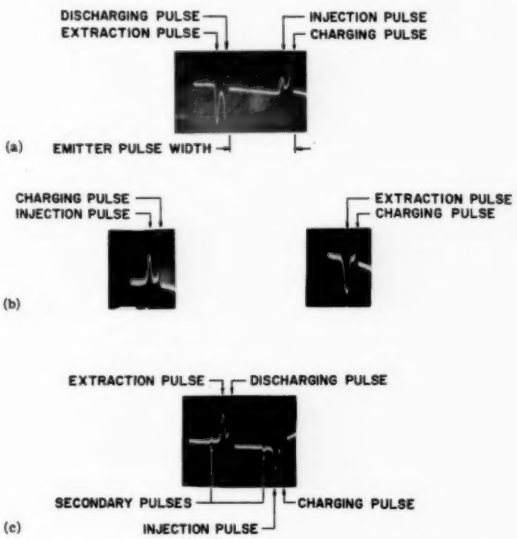

FIG. 2. Typical signals observed in the modified drift mobility experiment for a 35 ohm-cm, p-type, Ge sample. The ohmic sweeping pulse is balanced off. Time is increasing from right to left. The time of transit between the emitter and collector of the bulk pulses is about 4 microseconds. The charging and discharging pulses arise from the differentiation of the emitter pulse. (a) Typical oscilloscope trace for small negative pulse applied to emitter electrode. (b) Change of signal type arriving at collector electrode resulting from increase of voltage applied to emitter electrode. (c) Secondary pulses appearing later in time at collector electrode.

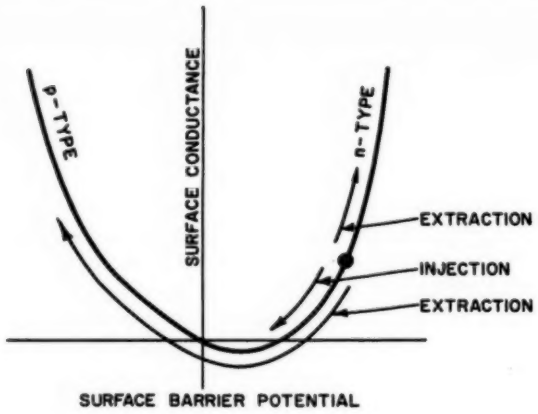

FIG. 3. Surface conductivity versus barrier potential diagram for a p-type sample.

tential diagram as shown in Fig. 3 in analyzing the observed signals. For a p-type sample with an n-type surface, extraction will occur if the surface is made more n-type and injection will occur if the surface is driven towards the minimum of the curve shown. If, however, the surface is driven strongly up the p-branch of the curve shown, extraction will be observed in the bulk.

This technique can be applied with advantage to the study of semiconductor surfaces. The interpretations of this experiment are more direct than in the usual field-effect experiment where it is necessary to analyze the measurements of sample conductivity. The following information about the surface can be obtained by means of this experiment:

1. The surface type is immediately ascertained from the polarity of the emitter pulse and the sign of the collected pulse. For example, in Fig. 2(a) the surface is n-type with respect to the bulk.

2. The surface barrier potential can be determined from observation of the changes in the collected pulse as the voltage of the emitter pulse is changed. The oscilloscope traces in Fig. 2(b) show how, with the proper pulse polarity, the minority carrier pulse in the bulk can be changed from injection to extraction by increasing the voltage pulse applied to the emitter electrode and thereby changing the surface from n-type to p-type with respect to the bulk, as is indicated in Fig. 3. As the voltage of the emitter pulse is increased, the collected signal at first increases but then decreases and even changes sign. The change in surface type begins to take place (i.e., when the signal begins to decrease) when the emitter voltage is 3.6. From this voltage and the measured capacitance, the charge induced in the condenser is calculated to be 2×10^{-9} coulomb. Because of the fast rise-time (10^{-8} sec) of the emitter pulse, it is assumed that this charge is associated entirely with the surface barrier, none being lost to surface states. Thus, the surface barrier potential is determined to be 0.05 volt from the curves of Kingston and Neustader.[5] This determination is in agreement with that measured from the change in contact potential with light for the same surface.

3. Information can be obtained about the relaxation effects in the surface by changing the repetition rate and the width of the emitter pulse as well as using a dc biasing field on the emitter electrode.

4. Information may perhaps be obtained about surface mobility. A well defined pulse [as shown in Fig. 2(c)] or step is often observed which appears later in time than the pulse corresponding to the drift of the carrier packet in the bulk. This second signal behaves as if it were associated with a mobility less than that corresponding to the bulk. From the numerous tests made, it is concluded that this second signal corresponds to the motion of a carrier disturbance in the surface. If this is so, the surface mobility can be measured directly.

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ACTIVATION ENERGY AND RELAXATION SPECTRUM OF DISLOCATIONS IN COPPER

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Several attempts have been made to evaluate the activation energy W and the frequency ω_A

Table I. Activation energy for dislocations in copper.

Authors	Frequency range kc/sec	W ev (molecule)$^{-1}$	Experimental data
Niblett and Wilks	0.4- 1.1	0.07	Niblett and Wilks
Niblett and Wilks	0.4-28.0	0.10	Niblett and Wilks; Bordoni
Donth	1.1-28.0	0.12	Niblett and Wilks; Bordoni
Caswell	...	0.15	Caswell
Paré	0.4-40.0	0.18	Paré; Caswell; Niblett and Wilks; Bordoni
Seeger, Donth, and Pfaff	...	0.11-0.13	Niblett and Wilks; Bordoni
Mason	0.4-28.0	0.09	Niblett and Wilks; Bordoni

for the low-temperature relaxation effect due to dislocations in copper. However, the measurements made on samples having the same impurity content and the same history cover a frequency range too narrow to make the evaluation of W accurate enough.[1] No better results are obtained when W is computed from measurements on samples having a different origin,[2] or when the theory of dislocation motion is employed to compute W by means of a single attenuation peak.[3] (See Table I.)

Some recent improvements in the experimental technique[4] made it possible to measure the attenuation from 1.8 kc/sec to 6.5 Mc/sec on the same samples of polycrystalline copper (grain size 6×10^{-2} mm) cross-rolled at high temperature and machined.

The frequency dependence of the temperature of maximum attenuation is shown in Fig. 1. All the experimental points lie very close to a straight line, and the relation between the frequency ω and the temperature T is given by

$$\omega=\omega_A e^{-W/kT}, \tag{1}$$

with

$$W=0.122\pm0.005 \text{ ev (molecule)}^{-1}, \tag{2a}$$

$$\omega_A=3.8\times(1.8)^{\pm1}\times10^{11} \text{ sec}^{-1}. \tag{2b}$$

The above values agree much better than any previous ones with the relation between W and ω_A derived by Seeger[5]:

$$\omega_A=\pi^2\frac{\sqrt{2}}{7}\frac{1}{b^4}\frac{W}{(\rho\mu)^{1/2}}, \tag{3}$$

where b = length of the Burgers vector $=2.55\times10^{-8}$ cm for copper, ρ = density = 8.94 g cm^{-3}, and μ = torsional modulus = 4.6 dyne cm^{-2}. From (2a) and (3) one finds $\omega_A=4.7\times10^{11}$, which is near enough[6] to (2b).

A relation slightly different from (1) has been recently obtained by Donth[7] from the theory of stochastic processes. However, the differences between the two equations are of the same order as the experimental errors. The ratio of the Peierls stress τ_p^0 to the torsional modulus μ can be computed from the experimental data by means of both theories,[8] with very similar results:

$$\tau_p^0\mu^{-1}=4.1\times10^{-4} \text{ (Seeger's theory)},$$
$$\tau_p^0\mu^{-1}=4.7\times10^{-4} \text{ (Donth's theory)}. \tag{4}$$

The measurements have also shown that W has the same value for all the dislocations, while there is a wide spectrum of values for ω_A. This may be seen from Fig. 2, where the experimental data obtained for the whole frequency range

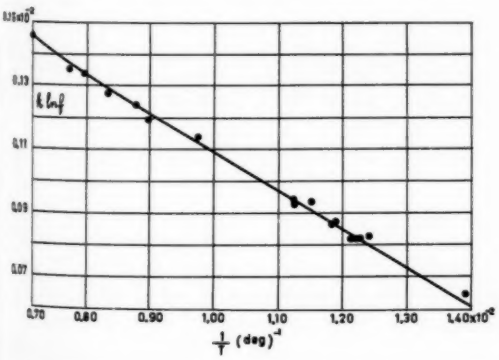

FIG. 1. Temperature dependence of the frequency of maximum attenuation.

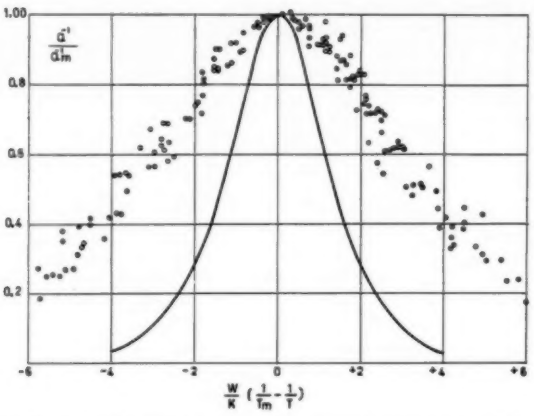

FIG. 2. Normalized attenuation values measured in the frequency range from 1.8 kc/sec to 6.5 Mc/sec (circles) compared with the attenuation curve for a single relaxation time (heavy line).

have been plotted against $Wk^{-1}(T_m^{-1}-T^{-1})$ after correction for background damping and normalization of the height of the peaks to unity. As all the points lie very close to a single line, W must be the same for all the dislocations. Moreover the line is broader than the curve for a single relaxation time; hence ω_A must be the center of a relaxation spectrum.

According to the above results the present theories are equivalent from a numerical standpoint, as far as the values of W and ω_A are concerned. However, both of them fail to explain the shape of the attenuation curve of Fig. 2, and give a better picture of the average dislocation behavior rather than of their actual motions. More detailed models seem therefore to be required to reconcile the existence of a frequency spectrum with that of a single value for the activation energy.

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[5]A. Seeger, Phil. Mag. 1, 651 (1956); see also W. P. Mason, Physical Acoustics and the Properties of Solids (D. Van Nostrand Company, Princeton, 1958), p. 271.

[6]W. P. Mason found the computed value twenty times larger than the experimental one; Paré found a ratio of 10^4 for the same quantities.

[7]See the papers quoted in references 2 and 3.

[8]For Donth's theory, the value of $\tau_p^0\mu^{-1}$ has been computed from a central point on the line of Fig. 1. However, the result depends only slightly upon frequency.

NUCLEAR MAGNETIC RESONANCE DETERMINATION OF ACTIVATION VOLUME FOR DIFFUSION IN LITHIUM*

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The purpose of this note is to call attention to the fact that the activation volume for self diffusion in solids can be measured by conventional nuclear magnetic resonance techniques. The activation volume for diffusion is important because it provides a measure of the extent to which the atoms of a solid relax inward on a vacancy, and because it may distinguish between vacancy and interstitial diffusion mechanisms. Specifically, the activation volume for diffusion is defined by[1]

$$V_{\text{act}} = (\partial \Delta F/\partial p)_T, \tag{1}$$

where ΔF is the free energy of activation and p the pressure. This can be written in terms of the diffusion coefficient D, the lattice spacing a, and the characteristic jump frequency ν_0, as

$$V_{\text{act}} = -RT\left[\frac{\partial \ln D}{\partial p} - \frac{\partial \ln \nu_0}{\partial p} - 2\frac{\partial \ln a}{\partial p}\right]. \tag{2}$$

Determination of the activation volume by tracer diffusion experiments requires measurement of the pressure dependence of D. To date, V_{act} has been determined for only a few elements, including the solid metals sodium[1] and lead.[2]

In the case of the lighter metals especially, the nuclear magnetic resonance line width undergoes a narrowing transition from its rigid lattice value

when the diffusion jump frequency becomes comparable to the resonance line width.[3] Alternatively, one may regard the spin-lattice relaxation time T_1 and spin-spin relaxation time T_2' as undergoing transitions. In particular, under conditions of diffusion,[4]

$$T_2' = 5cD/4\gamma^4\hbar^2 I(I+1)N, \qquad (3)$$

where c is the nearest neighbor distance, and N is the number density of resonant nuclei. Such line-narrowing transitions have been studied under normal pressure conditions for the metals lithium,[4] sodium,[4] aluminum,[5] and cadmium.[6]

Equation (3) may be used to eliminate D from (2):

$$V_{act} = -RT\left[\frac{\partial \ln T_2'}{\partial p} - \frac{\partial \ln \nu_0}{\partial p} - 6\frac{\partial \ln a}{\partial p}\right]. \qquad (4)$$

The addition to the third term comes from the volume dependence of c/N. This relation can be used to determine the activation volume for diffusion from the pressure dependence of the nuclear spin-spin relaxation time in the temperature range in which line narrowing due to diffusion occurs.

We have measured the line width $\delta\nu$ of the Li^7 resonance in lithium metal dispersed in mineral oil as a function of pressure to approximately 3000 atmos at several temperatures. The line width is quite strongly pressure dependent, roughly doubling at 3000 atmos, and within this pressure range $\ln\delta\nu$ is a linear function of the pressure. Figure 1 shows the results for two temperatures. In the case of lithium in the vicinity of room temperature, the spin-lattice relaxation contribution to the line width is negligible; T_2' may be taken as inversely proportional to $\delta\nu$. The contribution of the $\partial \ln a/\partial p$ term to V_{act} can be calculated from the isothermal compressibility; the $\partial \ln \nu_0/\partial p$ term can be shown to be negligible.[1] The values of

$$(\partial \ln T_2'/\partial p)_T = -(\partial \ln \delta\nu/\partial p)_T \qquad (5)$$

from the data of Fig. 1 then yield for the activation volume

$$V_{act} = 5.10 \text{ cm}^3 \text{ at } 3°\text{C},$$

and

$$V_{act} = 5.18 \text{ cm}^3 \text{ at } 27°\text{C}.$$

We estimate the error to be ± 0.25 cm^3 in each case. The ratio of the activation volume to the molar volume has the value 0.40, and should be compared with the values 0.52 for sodium[1] and 0.77 for lead.[2] If no relaxation of the atoms around a vacancy occurred, V_{act} would be at least as large as the molar volume.

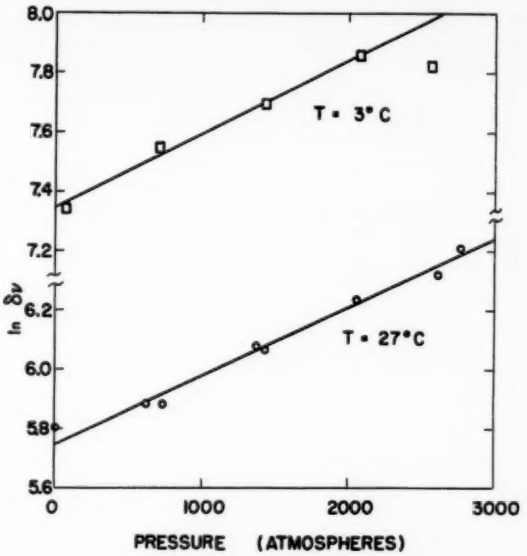

FIG. 1. Pressure dependence of the nuclear magnetic resonance line width in lithium metal. The departure of the last point in the 3°C data is due to freezing of the oil.

Tracer diffusion experiments are not feasible in the case of lithium so that magnetic resonance techniques provide the only access to its diffusion characteristics. Direct measurement of the pressure dependence of the relaxation times T_1 and T_2 by pulse techniques would yield somewhat more reliable values than the line width measurements described here. In addition, pulse methods would still be appropriate at higher temperatures where line width measurements become essentially unfeasible. For sodium and lithium it should be possible to study the temperature dependence of the activation volume in the solid phase and to observe the sharp decrease in V_{act} which must occur at the melting point.[2]

It is a pleasure to acknowledge helpful discussions with Dr. Wm. H. Jones, Jr.

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*Contribution No. 718. Work was performed in the Ames Laboratory of the U. S. Atomic Energy Commission.

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THEORY OF ELECTRON-PHONON INTERACTIONS

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The theory of the interaction of electrons and acoustic phonons in nonpolar crystals has been formulated in terms of a new set of basis states, whose wave functions are essentially Bloch functions that deform with the lattice. The major part of the interaction may then be calculated in terms of the strain tensor rather than the displacement of the lattice. A result of the theory is a generalization of the deformation potential theorem.[1]

The wave functions of the new basis states are

$$\eta_{n\vec{k}}(\vec{x},q)=\phi_{n\vec{k}}(\vec{y})\left[\frac{\partial(x_1,x_2,x_3)}{\partial(y_1,y_2,y_3)}\right]^{-1/2}, \tag{1}$$

where $\phi_{n\vec{k}}(\vec{x})$ is a Bloch function of wave vector $\vec{k}$ and band index n, where q stands for all the normal coordinates $q(\vec{f},s)$ of the lattice, and where $\vec{y}(\vec{x},q)$ is defined by

$$\vec{x}=\vec{y}+\vec{u}(\vec{y},q);$$
$$\vec{u}(\vec{y},q)=\sum_{\vec{f},s} q(\vec{f},s)\vec{v}(\vec{f},s)\exp(i\vec{f}\cdot\vec{y}). \tag{2}$$

Here s labels the lattice modes for a given wave vector $\vec{f}$, and $\vec{v}(\vec{f},s)$ is a polarization vector chosen so that $\vec{u}(\vec{R}_j^0,q)$ is the displacement of the center of mass of the jth unit cell from its equilibrium position $\vec{R}_j^0$. The expression in square brackets in (1) is the Jacobian of the transformation: this factor serves to make the $\eta_{n\vec{k}}(\vec{x})$ orthonormal. We call the $\eta_{n\vec{k}}$ the "orthogonalized deformed Bloch" (ODB) functions.

We have derived a Boltzmann equation for the diagonal elements of the density matrix in this ODB representation. The off-diagonal elements of the electron-plus-lattice Hamiltonian, H, in this representation, give the transition rates in the usual way.[2]

The matrix elements for ODB states have been calculated to first order in the $\partial u_\alpha/\partial y_\beta$, assuming an electron potential[3]

$$V(\vec{x},q)=V_0(\vec{y})+\sum_{\alpha\beta} U^{\alpha\beta}(\vec{y})\frac{\partial u_\alpha}{\partial y_\beta}, \tag{3}$$

in the vibrating crystal, where $V_0(\vec{x})$ is the potential in the undeformed crystal and $U^{\alpha\beta}(\vec{y})$ is an unspecified function with the lattice periodicity.

The Hamiltonian is conveniently split into three terms:

$$H=H^0+H^{\mathrm{I}}+H^{\mathrm{II}}.$$

H^0 is diagonal in the ODB representation, with matrix elements equal to the Bloch energy, $\epsilon_n(\vec{k})$, plus the energy eigenvalue of the harmonic lattice vibrations. H^{II} is a term associated with the fact that a moving lattice tends to drag the electrons with it, and is usually negligible.

H^{I} is the term of interest. The central result of this work is that

$$H^{\mathrm{I}}=\tfrac{1}{2}\sum_{\alpha\beta}\left(\mathfrak{D}^{\alpha\beta}\frac{\partial u_\alpha}{\partial y_\beta}+\frac{\partial u_\alpha}{\partial y_\beta}\mathfrak{D}^{\alpha\beta}\right), \tag{4}$$

where $\mathfrak{D}^{\alpha\beta}$, the deformation potential operator, is given by

$$\mathfrak{D}^{\alpha\beta}\equiv\frac{\hbar^2}{m}\frac{\partial^2}{\partial y_\alpha\partial y_\beta}+\frac{\hbar^2 k_\alpha}{im}\frac{\partial}{\partial y_\beta}+U^{\alpha\beta}(\vec{y}), \tag{5}$$

with $\vec{y}$ given by (2).

The matrix elements of $\mathfrak{D}^{\alpha\beta}$ are related to the deformation potential, $D^{\alpha\beta}(n,\vec{k})$, which as usual is defined by

$$\epsilon_n'(\vec{k})=\epsilon_n(\vec{k})+\sum_{\alpha\beta} D^{\alpha\beta}(n,\vec{k})\mathcal{C}_{\alpha\beta},$$

where $\epsilon_n'(\vec{k})$ refers to a crystal which has been homogeneously strained so that $R_\alpha^0\to R_\alpha^0+\sum_\beta\mathcal{C}_{\alpha\beta}R_\beta^0$. Using (3) and a change of variables similar to the first of Eqs. (2), one obtains[4]

$$D^{\alpha\beta}(n,\vec{k})=\mathfrak{D}_{n\vec{k};n\vec{k}}^{\alpha\beta}\equiv\int\phi_{n\vec{k}}^*(\vec{y})\mathfrak{D}^{\alpha\beta}\phi_{n\vec{k}}(\vec{y})d^3y. \tag{6}$$

The matrix elements[5] of H^{I} in the ODB representation [i.e., between the states $\eta_{n\vec{k}}(\vec{x})$ and

$\eta_{n'\vec{k}'}(\vec{x})$] are

$$\sum_{\alpha\beta}\tfrac{1}{2}\left[\mathfrak{D}_{n\vec{k};n'\vec{k}}{}^{\alpha\beta} I(n'\vec{k};n'\vec{k}') + \mathfrak{D}_{n\vec{k}';n'\vec{k}'}{}^{\alpha\beta} I(n\vec{k};n\vec{k}')\right] \times S_{\vec{k},\vec{k}'}{}^{\alpha\beta}, \quad (7)$$

where

$$I(n\vec{k};n'\vec{k}') \equiv \int w_{n\vec{k}}{}^{*}(\vec{y})\, w_{n'\vec{k}'}(\vec{y}) d^3y;$$

$$S_{\vec{k},\vec{k}'}{}^{\alpha\beta} \equiv \frac{1}{V}\int e^{-i\vec{k}\cdot\vec{y}} \frac{\partial u_\alpha}{\partial y_\beta} e^{i\vec{k}'\cdot\vec{y}} d^3y.$$

(V is the sample volume and $w_{n\vec{k}}$ is the periodic part of a Bloch function normalized in V.)

For intraband transitions, (7) reduces to

$$\sum_{\alpha\beta}\tfrac{1}{2}\left[D^{\alpha\beta}(n,\vec{k}) + D^{\alpha\beta}(n,\vec{k}')\right] I(n\vec{k};n\vec{k}') S_{\vec{k},\vec{k}'}{}^{\alpha\beta}. \quad (8)$$

The formula (8) for the matrix elements of H^{I} generalizes Bardeen and Shockley's result[1] to include transitions between arbitrary points in the zone, insofar as (3) is applicable for large phonon wave vectors, and (7) further extends it to include interband transitions.[6]

The author is indebted to P. J. Price for suggesting this research and for discussions and suggestions during its progress.

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[3] Equation (3) represents only that part of the change of the potential which is associated with the motion of the centers of mass of the unit cells, characteristic of acoustic modes. Moreover, it is not clear how valid it is for large phonon wave vectors. For pure long-wavelength acoustic modes it is evidently quite realistic. The special case $U^{\alpha\beta} = 0$ is the deformable ion model.

[4] Since k_α appears in (5) explicitly, the deformation potential $D^{\alpha\beta}(n,\vec{k})$ does not have the periodicity of the reciprocal lattice. However, $H^{I}+H^{II}$ does have this periodicity. Since $\mathfrak{D}^{\alpha\beta}$ is not symmetric in α and β, the interaction cannot be written in terms of the usual strain tensor $\frac{1}{2}(\partial u_\alpha/\partial y_\beta + \partial u_\beta/\partial y_\alpha)$.

[5] From this point on we leave out the lattice wave functions and integrals on q.

[6] We have not investigated the effect of spin-orbit coupling on this theory.

MICROSCOPY AT LIQUID HELIUM TEMPERATURES: PHASE TRANSITION IN SODIUM

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The pioneer work of Barrett[1,2] showed that at about 36°K, sodium begins to undergo a martensitic transformation from the body-centered cubic to the close-packed hexagonal phase. The analysis of his x-ray data suggested that at 4°K about 5% of the specimen had transformed to the hexagonal phase. Recent work by Martin[3] on the specific heat, and by Gugan and Dugdale[4] on the electrical resistance changes associated with the transition, leads one to believe that a much greater percentage of the metal must be transformed at 4°K. Tensile tests on sodium at 20° and 4°K in which a considerable hardening was observed[5] would also suggest a larger amount of transformation.

In order to investigate this problem further it was decided to examine the surface of a sodium specimen under high magnification while it was at a temperature near that of liquid helium. For this purpose a small portable cryostat was designed which would fit under the objective of an ordinary standard microscope. The sodium specimen which was in thermal contact with a bath of liquid helium was mounted about 1 mm from the top of the cryostat, and it was illuminated and viewed through a thin glass window. A copper-constantan thermocouple embedded in the sodium showed that its temperature was approximately 6°K. With this arrangement the specimen surface (which was cut, in situ, under vacuum at room temperature) could be examined using a standard 16-mm objective.

The observations show that the transformation is much greater than the 5% originally suggested. A typical photograph is shown in Fig. 1. Nearly the whole of the metal is covered by martensite plates which, due to the change in volume, are tilted relative to the surface thereby giving regions of high contrast. The amount of transformation is by no means homogeneous over the

FIG. 1. The surface of a sodium specimen at 6°K (magnification×200).

entire surface, but assuming that the area of contrast on the surface is proportional to the total metal which has transformed, then at 20°K about 45% is in the new phase. At 6°K no further observable transformation has occurred. This would agree with the results deduced from the electrical resistivity experiments of Gugan and Dugdale.[4] After the specimen is warmed up to 90°K very faint traces of the plates can be seen in some places, but in general the surface of the metal reverts to its original appearance.

Figure 1 should serve as a warning to those people who still think that sodium can be considered as an ideal metal at low temperatures.

This work which was done at the Clarendon Laboratory, Oxford, forms part of a joint program of research between the Laboratory and the Metallurgy Division of the Atomic Energy Research Establishment, Harwell.

A full description of the apparatus and the experiments will be published later.

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SPIN-PHONON INTERACTION IN RUBY

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In the "hot phonon" theories of paramagnetic relaxation,[1,2] saturation of the spin resonance absorption causes heating of those lattice modes whose frequencies lie in the neighborhood of the absorption line. As a consequence of this selective phonon heating it is to be expected that energy will be transferred from those spins which are in a saturation field, via the phonons, to any other spins in the crystal which have the same resonance frequencies, but which may not themselves be in the microwave saturation field. This mechanism has been proposed as a possible explanation for the uniform spin saturation in nonuniform saturating fields observed by Morris, Kyhl, and Strandberg.[3]

We have attempted to observe the spatial transfer of spin temperature in synthetic ruby crystals containing 0.05% and 0.005% Cr_2O_3 at 1.7°K and 20°K. The experiment consisted in observing the resonance absorption, in one end of a crystal, at a power level well below saturation as a function of the saturation power level applied to the other end. Dransfeld[4] has reported a similar experiment on ruby with 1% Cr_2O_3, but at considerably lower saturation spin temperatures.

The experimental geometry is shown in Fig. 1. The ruby crystal, in the form of a circular cylindrical rod 0.1 inch in diameter and 1.5 inches long, was supported at one end and extended between two cavity resonators. The rectangular cavities operated in the TE_{012} mode and were tuned to the same frequency, 9.3 kMc/sec. The saturation power could be applied to either cavity, and the low-power absorption line observed in the other. In order to reduce the direct electromagnetic coupling between the resonators they were spaced $\frac{1}{4}$ inch apart, and were oriented such that the microwave fields at the crystal were orthogonal. In the experiments at 20°K the cavities were pressurized with He gas, and at 1.7°K they were filled with liquid He.

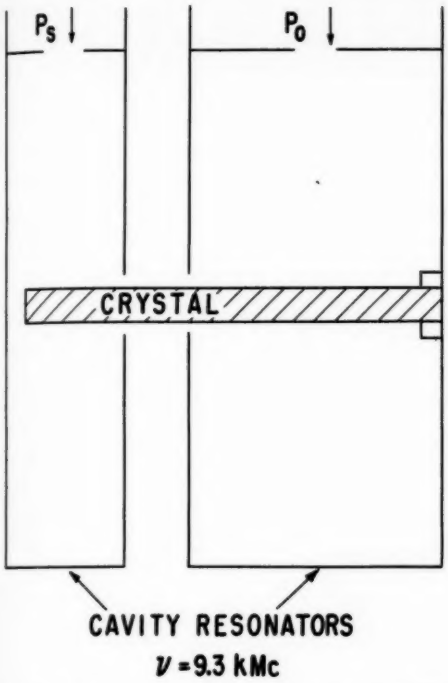

FIG. 1. Schematic drawing of the dual cavity system, shown with saturating power (P_S) in the left-hand resonator. The observation power $P_0 \ll P_S$.

Since there was a small amount of direct coupling present, this had to be subtracted in the proper phase from the observed signal. After making the subtraction, the detection sensitivity in the observation cavity, expressed as a minimum detectable change in spin temperature, was 0.2°K at 1.7°K, and 2°K at 20°K.

The maximum spin temperatures achieved in the saturation cavity were 30 000°K and 100 000°K, for the experiments at 1.7°K and 20°K, respectively. No detectable change in spin temperature was observed in the other cavity. Dransfeld's results, obtained for a saturation spin temperature >30°K, are thus consistent with ours.

On the basis of the hot-phonon theory it is possible to estimate an expected increase in the spin temperature at the observation cavity. We consider the following one-dimensional heat flow problem in the crystal. The heat input to the crystal is the total microwave power absorbed by the spins in the saturation cavity, and is distributed uniformly over that portion of the crystal. The surface cooling rate is assumed proportional to the difference in temperature between the heated lattice modes and the bath, and the boundary conditions on the end faces are that the heat conducted to them is equal to the heat lost from them by surface cooling. Our principal interest is in whether our experimental result is compatible with the prediction of the hot-phonon theory that it is the lattice-bath relaxation, and not the spin-lattice relaxation, which limits the total relaxation rate.[5] We therefore assume that the spin-lattice rate is very fast and the measured relaxation time is all lattice-bath. The surface cooling constant can then be evaluated by solving the time-dependent problem, taking the steady-state solution (with unknown cooling constant) as initial condition, and equating the decay rate to the measured relaxation time. In order to obtain a lower limit on the expected temperature rise in the observation cavity, we assume that the heat input spreads rapidly over the whole lattice spectrum. For the experiment at 1.7°K on the 0.05% Cr_2O_3 crystal, the steady-state solution (with cooling constant evaluated as above) gives an expected temperature rise in the observation cavity of 2.5°K. A smaller band of lattice modes would give a larger increase.

Since our minimum detectable temperature is more than an order of magnitude smaller than that expected from the hot phonon theory it must be concluded that, at least in ruby, it is the spin-lattice relaxation and not the lattice-bath time which limits the total rate, and thus the hot-phonon theories are not applicable in this material.

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MONOENERGETIC POSITRONS IN THE DECAY OF Bi^{206}.

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It was pointed out by Sliv[1] that a competing process to the ordinary internal pair formation

must exist in which monoenergetic positrons are produced. This process can be expected after a K-capture into an excited nuclear state from which a gamma ray is emitted with an energy greater than $2mc^2$. In such a case, beside the usual decay modes, such as gamma-ray emission, internal conversion, and internal pair formation, another de-excitation process is possible. After the K-capture a vacancy exists in the K-shell. An electron-positron pair can be created, the electron captured in the K-shell, and only the positron emitted. The positron energy has a well-defined value $E_K{}^{e+} = E_\gamma - 2mc^2 + E_K$, where E_γ is the decay energy and E_K is the binding energy of a K-electron.

To search for this effect we have chosen the electron capture decay of Bi^{206} into the 3.403-Mev level of Pb^{206}. [The activity was produced by the reaction Pb^{206}(15-Mev $d, 2n$)Bi^{206} in the cyclotron at Birmingham.] The 3.403-Mev level decays predominantly by electric dipole radiation of 1.72 Mev. Measuring the positron spectrum with a lens spectrometer, we have found a positron line (see Fig. 1) which we interpret as the monoenergetic positrons from the 1.72-Mev transition. The energy is about 90 kev higher than the end point of the internal pair spectrum of the same transition. The form of the line as drawn in the figure corresponds to a least-squares fit of the line profile of our apparatus. This profile was measured with a conversion line of the investigated source. A check of the electron-positron separation of the spectrometer excludes conversion electrons as a possible cause of the observed line.

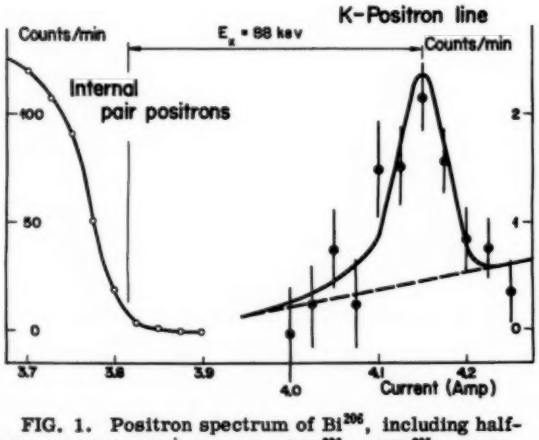

FIG. 1. Positron spectrum of Bi^{206}, including half-life correction. β^+-spectra of Bi^{206} and Bi^{205}, which were about 10 times the maximum counting rate of the positron line, have already been subtracted.

We compare the positron line with the K-conversion line of the 803-kev transition. The intensity ratio of these two lines is $N_{K1720}{}^{e+}/N_{K803} = (0.98 \pm 0.17) \times 10^{-5}$ (rms error). From this we determine the emitted number of positrons per gamma quantum of 1.72 Mev. We make use of the gamma-ray intensities of the 803-kev and the 1720-kev transition, measured by Novakov et al.[2] together with the theoretical K-conversion coefficient of the 803-kev $E2$ transition. The number of positrons per gamma quantum turns out to be $N_{K1720}{}^{e+}/N_{\gamma 1720} = 2.61 \times 10^{-7}$. This information allows us to compute the lifetime of the 3.403-Mev level.

Obviously the intensity of the positron line is proportional to the probability that the hole in the K-shell is not filled again by an electron of an outer shell before the de-excitation of the nuclear level. In the case where the mean life of the K-vacancy, τ_K, is long compared with the mean life of the nuclear level, $N_K{}^{e+}/N_\gamma$ is given by the "conversion coefficient for the formation of monoenergetic positrons," i.e., the probability of the emission of a monoenergetic positron to the probability of a gamma-ray emission, calculated by Sliv.[1] In order to get $N_K{}^{e+}/N_\gamma$ in the present case, where $\tau_K \ll \tau$, the conversion coefficient of Sliv has to be multiplied by the probability that the hole is not filled again before the de-excitation of the nuclear state. Because this probability is simply τ_K/τ, we can calculate the nuclear mean life τ in terms of the lifetime of the K-vacancy.

Before doing this one has to correct for the fact that the 3.403-Mev level is partly reached by L-capture. If we denote by P_K the probability that the 3.403-Mev level is excited by K-capture, we get the following relation:

$$N_{K1720}{}^{e+}/N_{\gamma 1720} = P_K \alpha_{K1720}{}^{e+} (\tau_K/\tau), \quad (1)$$

where $\alpha_K{}^{e+}$ is the conversion coefficient of Sliv. P_K can be calculated from the energy difference between the ground state of Bi^{206} and the 3.403-Mev level. We have found a positron spectrum in the decy of Bi^{206} with an intensity of about 4×10^{-5} per disintegration (maximum energy 1000 ± 60 kev; $\log ft = 10.1_{-0.3}{}^{+0.2}$). If we suppose that this transition leads to the 1.684-Mev level (suggested by the estimated energy difference between the Bi^{206} and Pb^{206} ground states[3,4]) the capture decay energy is fixed and we compute $P_K = 0.77$.[5] From Eq. (1) we get $\tau = 152\,\tau_K$, taking $\alpha_{K1720}{}^{e+} = 5.16 \times 10^{-6}$ from the work of Sliv.

The mean life τ_K of the K-hole is given by the

level width of the K-level, which can be determined from the line width of x-rays. We got 57 ev for the K-level width of Pb, extrapolating the most accurate value for Bi proposed by Mladjenović[8] (Z-dependence according to a Z^4 law). With this value, the mean life of the 3.403-Mev level of Pb^{206} is $\tau = 1.8 \times 10^{-15}$ sec. Compared with the transition probability of a single proton,[7] the 1.72-Mev electric dipole transition is slowed down by a factor of about 50.

We are deeply indebted to Mr. F. Stuber for the chemical preparation of the source. We would like to thank the cyclotron group of the University of Birmingham for the irradiations.

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PRODUCTION OF POLARIZED PROTON BEAMS*

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Several techniques[1] have been proposed to provide beams of polarized protons. One of the earliest proposals was put forward in a paper by Lamb and Retherford[2] where a basic outline was given for a method utilizing the Lamb shift. This technique had been considered to be unfeasible because of the problems of obtaining intense beams of metastable hydrogen atoms.[3]

In this note a method is proposed to produce relatively intense beams of metastable hydrogen atoms and their subsequent ionization. In addition, preliminary experiments related here indicate the feasibility of the proposed technique.

A proton beam can be formed by conventional methods with a beam energy of the order of 10 kev. The metastable hydrogen atoms are then produced by passing this proton beam through a chamber of hydrogen gas, or another appropriate gas at pressures of the order of 10^{-4} to 10^{-3} mm Hg. The high cross sections[4] (of the order of $0.5\pi a_0^2$) for pickup (or charge exchange) of an electron by a proton in the beam provides an efficient mechanism for the production of metastables.

After one removes the residual charged particles by deflection, the neutral beam, consisting of metastables and ground-state atoms, is then passed into a magnetic field of the order of 575 gauss. In the magnetic field the lower component of the $2S$ state becomes degenerate with a $2P$ level and decays to the ground state. Since the upper $2S$ state also possesses a hyperfine splitting, an appropriate rf field enables one to quench out the desired hyperfine component. The resultant neutral beam now consists of metastable atoms having polarized nuclei and ground-state atoms and molecules.

The metastable state can be preferentially ionized by the introduction of ultraviolet radiation in the region of 3650 A using a high-pressure mercury arc source. At this point the emerging charged particles consist of polarized protons. These can be separated from the neutral beam into an accelerator by an appropriate deflecting field.

Preliminary experiments have been performed to test this proposal for the method of producing a beam of metastable hydrogen atoms. In these experiments a 10-kev beam was used from a standard radio-frequency-excited ion source. A diaphragm having an aperture of 1-mm diameter, 50 cm from the ion source, defined a low-intensity beam. This diaphragm also served to provide a back pressure of hydrogen gas at 2×10^{-4} mm Hg in the volume nearest the ion source. Although the mean free path at these pressures was many times the path length in the gas volume, a reasonable amount of pickup into the $2S$ state was detected. Both magnetic and electric fields were provided at the point where the beam of neutral and charged particles emerged from the diaphragm (the prequenching region). These fields, which could be applied independently, served to deflect the charged particles from the beam, as well as to quench the neutral metastable beam. An adjustable electrostatic or magnetic field was placed 50 cm from the diaphragm behind a second aperture. These second fields provided a quenching region for the detection of the metastable beam. A small region in the center of the volume

containing the second quenching fields was viewed by a Lyman alpha Geiger counter.[5] The purpose of the experiment was to measure the decay of the atoms from the 2S state as a function of the intensity of the electrostatic and magnetic quenching fields at the position of the detector and in the prequenching region.

Since the atoms were in the quenching field before they passed the field of view of the Lyman alpha counter, one would expect the counting rates to reach a maximum at a given field strength and then to drop off as the field is increased further, since at high field intensity the metastables would be quenched before they could reach the region viewed by the counters.

This behavior was observed for both electric and magnetic quenching. The effective quenching rates are of course different, since the magnetic quenching field not only provides a motional electric field perturbation, but also changes the energy of the states. The magnetic quench curve implies that electric quenching does not produce additional excited states giving rise to Lyman alpha radiation. The results given in Fig. 1 show a typical electrostatic quenching curve and the field dependence seems to agree with expectation. The qualitative agreement of the theoretical and experimental curves supports the evidence that the electric fields are not giving rise to extraneous excited states which would then provide Lyman alpha radiation.

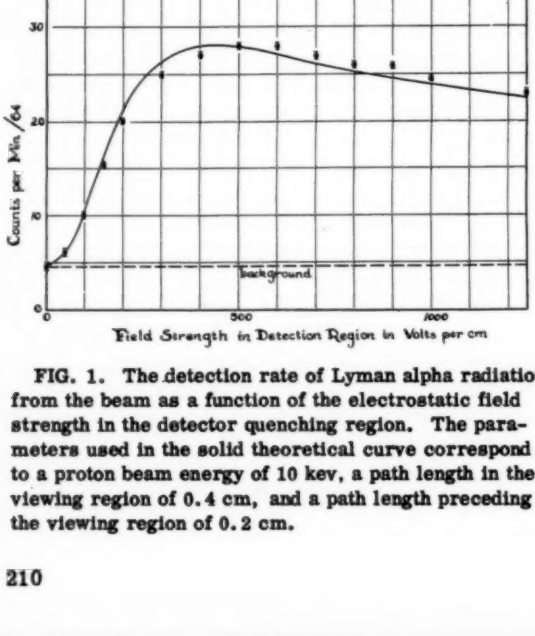

FIG. 1. The detection rate of Lyman alpha radiation from the beam as a function of the electrostatic field strength in the detector quenching region. The parameters used in the solid theoretical curve correspond to a proton beam energy of 10 kev, a path length in the viewing region of 0.4 cm, and a path length preceding the viewing region of 0.2 cm.

One can set the detector quenching fields so that the counter records the maximum rate, and subsequently observe the effect on this rate due to the field intensity in the prequenching region. Figure 2 shows the effect of prequenching electrostatically, where an electric field was used in the detector quenching region. It is important to notice that the theoretical predictions are very sensitive to the precise field configurations, the viewing region, and the beam velocity. The comparison of theory to experiment therefore does not represent high precision, but rather the fact that the experiments behave correctly in a qualitative way. A rough estimate of Lyman alpha intensity indicates that the pickup cross section is $>0.01\pi a_0^2$.

The results of the preliminary experiments seem to indicate that a beam of metastable atoms can be achieved by the pickup process. The ionization of these metastables will provide a 50% polarized beam without the introduction of radio-frequency quenching. It is of interest to notice

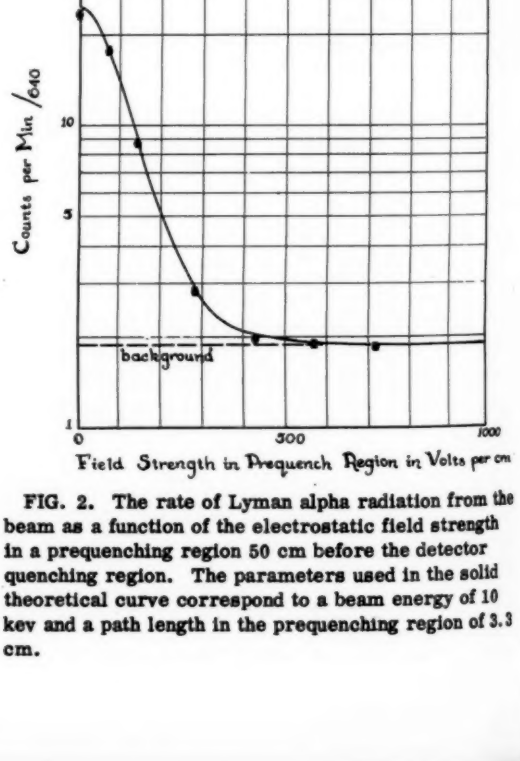

FIG. 2. The rate of Lyman alpha radiation from the beam as a function of the electrostatic field strength in a prequenching region 50 cm before the detector quenching region. The parameters used in the solid theoretical curve correspond to a beam energy of 10 kev and a path length in the prequenching region of 3.3 cm.

that the photoelectrons from the metastables are also polarized, and in principle could be collected to form polarized electron beams. Further experiments on the production of metastables, photoionization, and acceleration are now in progress in cooperation with N. Heydenburg, G. Temmer, and J. Weinman at the Department of Terrestrial Magnetism.

The authors wish to thank Dr. Donald E. Kerr for many helpful suggestions. They are also indebted to Dr. N. Heydenburg, Dr. G. Temmer, and Dr. J. Weinman for valuable discussions concerning the problem of producing polarized beams in an accelerator.

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SEARCH FOR THE DECAY $\mu^+ \to e^+ + \gamma$*

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(Received February 5, 1959)

It has been pointed out recently that the existence of a heavy charged boson with the properties required to give couplings consistent with the universal V-A Fermi interaction will lead to the occurrence of the decay $\mu^+ \to e^+ + \gamma$.[1,2] The branching ratio $\rho = R(\mu \to e + \gamma)/R(\mu \to e + \nu + \bar{\nu})$, when calculated by assuming such an intermediate meson coupling, is of order 10^{-4}. Previous experiments have given a value of $\rho \leq 2 \times 10^{-5}$.[3] This note is to report the results of an experiment in progress at the University of Rochester synchrocyclotron, designed to obtain a more precise value of ρ.

The arrangement of the detection apparatus is shown in Fig. 1. A 32-Mev π^+-meson beam is stopped in the carbon target. The μ^+-mesons arising from π^+ decay also stop and decay in the target. The stopped π^+ beam had an intensity of 5×10^5 mesons/minute. Two counter telescopes are placed so as to observe coincident positrons and gamma radiations emitted in opposite directions from the target.

The positron telescope consists of two scintillation counters e_1 and e_2 followed by a water Čerenkov detector, e_3 (see Fig. 1). Only positrons with energy $E \geq 35$ Mev were counted by this telescope. The solid angle subtended by this counter array at the source was 5% of 4π.

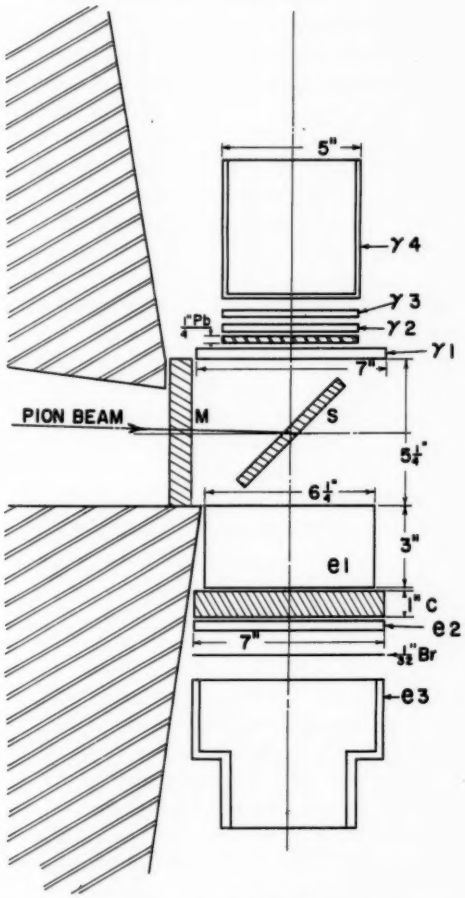


FIG. 1. Arrangement of detection apparatus with respect to the 32-Mev pion beam. M is the meson moderator and S is the target in which the pion beam stops. One inch of carbon absorber is placed between e_1 and e_2, 1/32 in. of brass between e_2 and e_3, and 1/16 in. brass between γ_2 and γ_3 (not shown).

For calibrating purposes, the positron telescope was first placed in a 95-Mev/c electron beam, and the delay and gain of each of the counters adjusted to form coincidences. Checks of these settings were made by observing coincidences in the electron beam moderated to 50 Mev/c, and also by observing μ^+-decay positrons with $E > 35$ Mev/c. The counting rate observed for μ^+-decay positrons is consistent with that expected from the solid angle and approximate estimates of scattering and radiation losses. The intrinsic efficiency of the positron telescope was estimated to be 75% for 55-Mev positrons.

The γ-ray telescope consists of three scintillation counters γ_1, γ_2, γ_3, and a water Čerenkov counter γ_4; γ_1 is an anticoincidence counter. A $\frac{1}{4}$-in. Pb radiator is placed between γ_1 and γ_2.

The γ-ray telescope was also calibrated in the 95-Mev/c electron beam; the gains and delays were adjusted first for $\gamma_2\gamma_3\gamma_4$ coincidences and then for $\bar{\gamma}_1\gamma_2\gamma_3\gamma_4$ events. The anticoincidence efficiency was about 99%.

A further calibration was carried out with high-energy photons from π^- capture in a liquid hydrogen target. The effective solid angle and intrinsic efficiency of the γ-telescope were measured by observing the fraction of π^- mesons absorbed in hydrogen that gave γ counts at several distances of the telescope from the hydrogen target. The measured solid angle in the final experimental arrangement is about 4.5% of 4π. The efficiency for 55-Mev γ-rays is estimated to be 19%, by taking it as 75% of the measured efficiency for the somewhat higher average energy gamma-ray spectrum from π^--capture. The minimum electron energy for detection was about 12 Mev.

A block diagram of the electronics is given in Fig. 2. To reduce the random e-γ coincidence rate, a "fast-slow" coincidence circuit was used that took advantage of the speed of the 6655A multipliers to reduce the effective resolving time for e-γ coincidences.

The relative delay of the two telescopes was adjusted by aligning both telescopes in the 95-Mev electron beam, which has enough range to penetrate both.

Two runs have been completed to date. The first run, on $3\text{-}4\times10^8$ stopped mesons taken with a relatively high background random rate, gave 208 total counts (real plus random), 170 random counts, or

$$\rho = (1.7 \pm 0.9)\times10^{-5}.$$

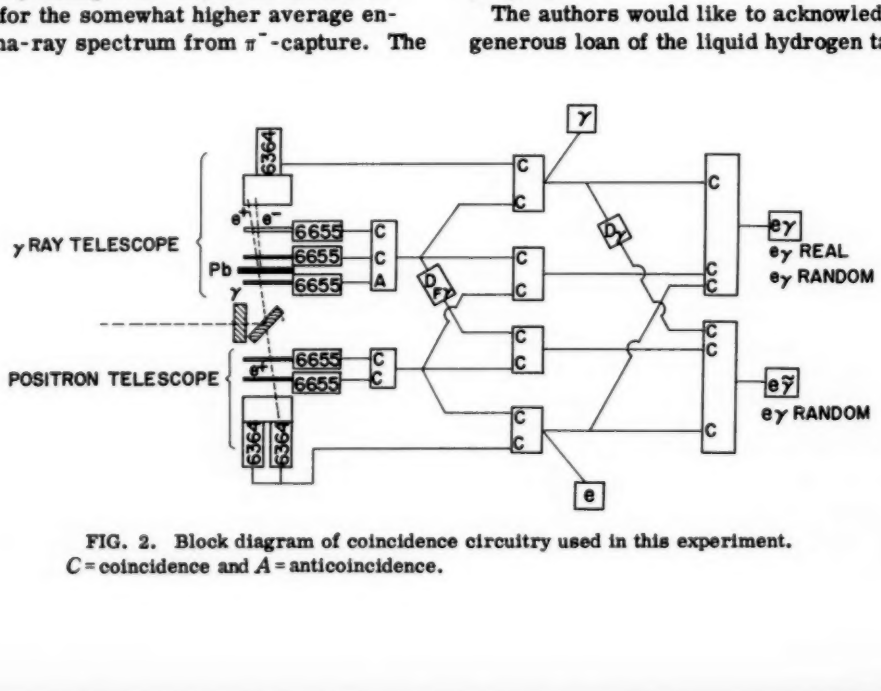
The second run, at a lower background rate, gave for 5.3×10^8 mesons 240 total, 237 random counts, or

$$\rho = (0.9 \pm 7)\times10^{-6}.$$

The quoted errors are statistical standard deviations.

The possible systematic errors of the experiment are such as to permit absolute estimates to be in error as much as 25 to 50%. Accordingly, we feel that the present experiments indicate that the probability that the branching ratio exceeds 1.0×10^{-5} is not over 0.5. Further work is in progress.

The authors would like to acknowledge the generous loan of the liquid hydrogen target by Dr.


γ
6364
e⁺
e⁻
γ RAY TELESCOPE
6655
6655
Pb
6655
γ
C
C
A
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e⁺
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6364
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Sidney W. Barnes, and the helpful assistance of Mr. Lawrence Ota.

*This work was supported in part by the U. S. Atomic Energy Commission.

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PRECISE DETERMINATION OF THE MUON MAGNETIC MOMENT*

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(Received February 5, 1959)

When the existence of polarized meson beams and the anisotropic electron emission in decay was first detected,[1] it became possible to measure accurately the magnetic moment of the μ meson. This was first accomplished by a precession technique[1,2] to an accuracy of 0.7%, and to 0.06% by the resonance method,[3] which uses a radio-frequency magnetic field to induce transitions of the muon spin from a state of alignment to one of antialignment (or vice versa). The latter technique does not lend itself to use for experiments of improved accuracy at higher frequencies because of the large rf field required. A stroboscopic method was adopted,[4] in which the muon is brought to rest with its spin perpendicular to a magnetic field, H. The counting rate of the decay electrons in a particular direction varies in time as $\exp(-t/\tau)(1+a\cos\omega_H t)$, where a is the usual experimental electron asymmetry, τ the muon mean life, and ω_H, the rate of muon precession, is equal to $geH/2mc$, where m is the muon mass, e the electronic charge, c the speed of light, and g the muon g-factor. By use of a higher precession frequency and by a different method of utilizing time information of the muon decay, the experiment described in this Letter achieved a muon moment accuracy of 0.007%.

Early experiments[2-5] displayed directly on a pulse-height analyzer (using a time-to-pulse-height converter) the electron counting rate vs time. More accuracy (which means higher fields since the frequency uncertainty is fixed by the muon mean life) requires some method of folding the sinusoidally modulated decay curve into the memory of a pulse-height analyzer and of achieving better linearity than is possible with an analog time-to-pulse-height converter. The natural approach is to use a reference oscillator as in the resonance method, or as one input to a coincidence circuit whose other input is the electron count.[4] In the latter case a resonance curve is generated by the coincidence rate as the field or reference oscillator frequency is varied, although serious attention must be paid to reduce the systematic errors below the possible accuracy of this technique.

In this experiment a cw oscillator was used to measure both the muon and decay-electron times, modulo the period of the reference oscillator (86.200 Mc/sec). That is, one measures both the muon and electron phase with respect to this cw stable oscillator. The phase difference is generated electronically and stored in one of 10 channels of the pulse-height analyzer. If the reference frequency, ω, were exactly equal to the muon spin precession frequency, then the distribution-in-phase $[n(\phi)]$ for electrons emitted during a single period (of the 86.2-Mc/sec oscillator) occurring at delay time T, is independent of T. For $\omega_H \neq \omega$, $n(\phi)$ becomes $n(\phi+\alpha)$, i.e., retains its functional form but is shifted along the ϕ axis by an amount $\alpha = (\omega_H - \omega)T$. Thus a measurement of α vs T gives $\omega_H - \omega$ and thus ω_H itself.

The electrons are grouped into two time intervals. Those immediately following the stopping of the muon are referred to as the early electrons. The remainder are referred to as the late electrons. The average $n(\phi)$ for these two groups are displayed separately, and α determined for each. The quantity $(\alpha_{\text{early}} - \alpha_{\text{late}})$ is then a direct measure of $\omega_H - \omega$. This subtraction procedure has the important advantage of canceling out systematic errors due to uncertainties in the starting phases.

Figure 1 shows the counter, absorber, and target arrangement. An incoming μ meson is defined by a $123\bar{4}$ coincidence which serves to open a 2×10^{-7}-sec gate for the fast counter pulse in the μ timing channel and a 6-μsec gate that permits a 2×10^{-7}-sec gate to be opened by a decay electron in the electron channel. A forward emitted electron is defined by a $\bar{1}345$ coincidence, a backward one by a $23\bar{1}\bar{4}$ coincidence. Either of

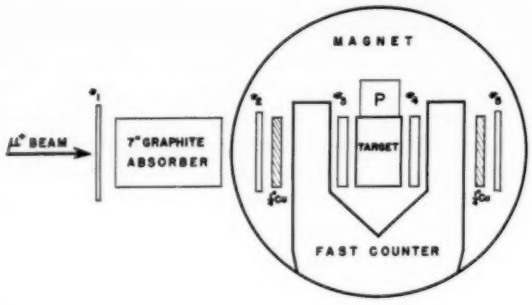

FIG. 1. Experimental arrangement. P=nuclear magnetic resonance probe.

these events opens the electron timing channel but the information obtained is routed to different places in final storage.

A block diagram of the time-analyzing equipment is shown in Fig. 2. The phases of the fast counter pulses for μ's and electrons relative to the free-running 86-Mc/sec source are automatically determined by beating pulsed 85-Mc/sec oscillators against the 86-Mc/sec signal and measuring the starting phase of the beat note. The 85-Mc/sec oscillators are triggered by overclipped fast counter pulses put through a zero-crossing detector to obtain a time independent of pulse height. Our timing accuracy is better than $\pm 5\times 10^{-10}$ second. Suitable delays and gates are included to remove the effects of starting transients. The pulsed oscillators are used only for phase determination so that their phase and frequency stability are not important. A condenser ramp is charged with the μ phase and discharged with that of the electron, thus obtaining an output phase whose height is proportional to $\phi_\mu - \phi_e$. The output is stored in a Penco 100-channel pulse-height analyzer which had been modified so that it was used as four 10-channel analyzers. The output pulse was addressed to the appropriate group of channels depending on whether it was emitted in the forward or backward direction and whether it occurred early or late in the gate interval.

Measurements were performed at several frequencies on either side of resonance. A plot of $\alpha_{\text{early}} - \alpha_{\text{late}}$ was obtained as a function of the magnetic field in which the μ precesses. This is a nearly linear curve whose intercept, $\alpha = 0$, corresponds to the magnetic field at which $\omega_\mu = \omega$. This obviates the necessity of knowing accurately the timing intervals for the early and late electrons. The results of our measurements are given in Table I in terms of the ratio of the

Table I. Ratio of the resonant frequency of the μ meson to that of the proton in water in the same magnetic field.

Target	Measured f_μ/f_p	After correction for Knight shift
Aluminum	3.1850±0.0002	3.1847±0.0003
Copper	3.1850±0.0002	3.1848±0.0003
Bromoform	3.1846±0.0002	3.1846±0.0002

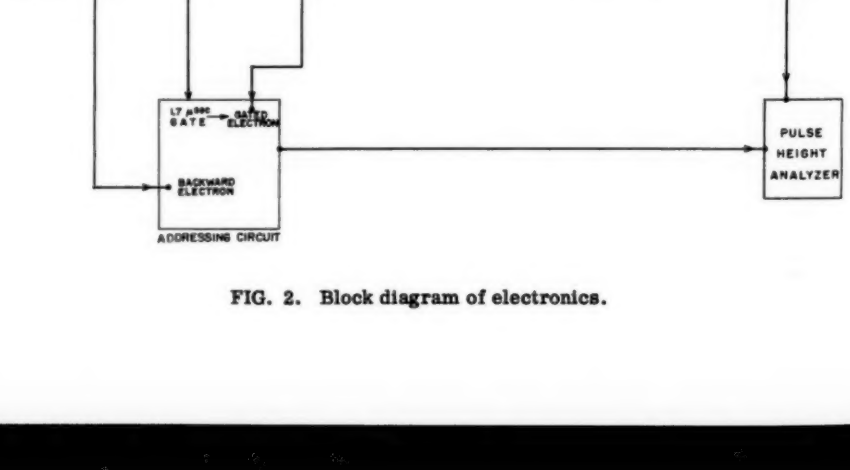

FIG. 2. Block diagram of electronics.

resonant frequency of the μ meson to that of the proton in water in the same magnetic field. The results for copper and aluminum are corrected by a rough calculation of the Knight shift[6] and then agree with that for $CHBr_3$ in which there should be no diamagnetic shift to the accuracy quoted. It should be noted that our result for aluminum disagrees with that of the Chicago group[4,7] to about twice their stated error.

Using our result for $CHBr_3$ and the value for the μ-meson mass given by Crowe,[8] the g-factor is found to be $2(1.0020 \pm 0.0005)$ as compared to $2(1.00116)$ as predicted by theory. Of even greater interest is the lower limit for the μ-meson mass given by Crowe from a determination of μ-mesonic x-rays[9] of $(206.77 \pm 0.04) m_e$. This yields $g > 2(1.00154 \pm 0.00022)$ which is also in disagreement with theory. It should be noted that the error arises mostly from the measurement of the mass which stems from uncertainty in the calculation of the vacuum polarization correction.

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†Also at IBM Watson Laboratory, Columbia University, New York, New York.

‡John Tyndall Fellow.

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NEUTRAL CASCADE HYPERON EVENT*

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The existence of a neutral cascade hyperon Ξ^0 has been predicted theoretically,[1] on the basis of the strangeness theory of Gell-Mann and Nishijima, as the neutral counterpart of the negative cascade hyperon,[2] Ξ^-, which decays by $\Xi^- \to \pi^- + \Lambda$.

In an attempt to establish the existence of this particle the Lawrence Radiation Laboratory 15-inch hydrogen bubble chamber was operated in a separated beam of (1.15 ± 0.02)-Bev/c K^- mesons produced by the Bevatron. Two Cork-Wenzel-Lambertson parallel-plate spectrometers[3] were used to remove pions from the beam. Typical operating conditions gave ≈ 1.5 K^-, ≈ 0.2 π^-, and ≈ 4.5 beam μ^- mesons per picture.[4] The total number of K^- mesons through the chamber was about 10^5.

A large number of K^- interactions in hydrogen were observed; among them were some 500 single V^0 events, resulting from the reactions

$$K^- + p \to \overline{K}^0 + n, \tag{1a}$$

$$K^- + p \to \Lambda + \pi^0, \tag{1b}$$

$$K^- + p \to \Sigma^0 + \pi^0. \tag{1c}$$

In any of these, additional π^0 mesons may also have been produced.

On the other hand, only seven double V^0 events were observed. Since the reactions (1) lead only to single V^0's, whereas associated production by π^- mesons leads to double V^0's in about 20% of the interactions, the strikingly small ratio of double V^0 events to single V^0 events again shows that we are dealing principally with K^- interactions.

Six of the double V^0's were clear cases of associated production by π^-, five being

$$\pi^- + p \to \Lambda + K^0, \tag{2a}$$

and one

$$\pi^- + p \to \Sigma^0 + K^0, \quad \Sigma^0 \to \Lambda + \gamma. \tag{2b}$$

Most of these were produced by pions of some-

what less than the K^--beam momentum.

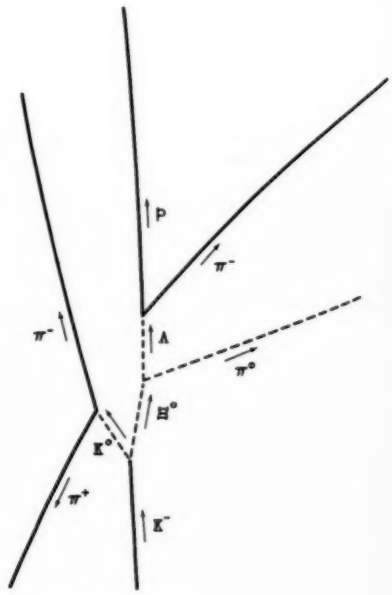
The remaining event is the one being reported here. A photograph and a diagram giving our interpretation of the event are shown in Fig. 1. The angles and momenta of the left-hand V^0 are consistent with K_1^0 decay, and are inconsistent with Λ decay. The K^0 momentum and angle of emission are consistent with the reaction $\pi^- + p \rightarrow \Sigma^0 + K^0$ of a beam-momentum pion.

The two charged tracks of the right-hand V^0 are consistent with Λ decay, giving $Q = 37.2 \pm 2.7$ Mev (accepted value 37.4 Mev). However, the decay is noncoplanar; i.e., the line connecting the end of the beam track and the vertex of the Λ fails by $7.0 \pm 0.7°$ (see Fig. 2) to lie in the Λ decay plane. This line also fails to lie in the production plane defined by the K^0 path and the beam track by $2.5 \pm 0.7°$. The latter discrepancy could be explained easily if the process were (2b), but to explain the lack of coplanarity of the Λ decay, using only well-established processes, we must invoke either (a) reaction (2b) followed by a β decay of the Λ, or (b) a scattering of the Λ in the hydrogen, or (c) an accidental coincidence of a K^0-meson production event with an unassociated Λ from the bubble chamber wall.[5]

Possibility (a) may be ruled out on kinematic grounds alone. Because of the large unbalance of transverse momentum, the electron and neutrino need more energy than would be available to them. The decay[6] $\Lambda \rightarrow p + e^- + \nu$, for the most favorable Λ momentum, fails to balance energetically by 48 Mev, or 3.7 standard deviations; the error is mostly in angle measurements. For such large discrepancies, angle errors do not have Gaussian distributions, and this large a discrepancy is not possible. A decay via $\Lambda \rightarrow p + \mu^- + \nu$ fits even less well; radiative decay, $\Lambda \rightarrow p + \pi^- + \gamma$,[7] may also be ruled out by similar arguments.

The second possibility, a Λ scattering, is likewise unsatisfactory. Choosing that initial Λ direction of motion for which the scattering angle would be smallest, one asks what the proton recoil range would be to account for the observed Λ. This turns out to be 4 mm, which would be clearly visible. To have a proton range small enough that there would be some doubt, namely

(a)



(b)

FIG. 1. Photograph and sketch of Ξ^0 event.

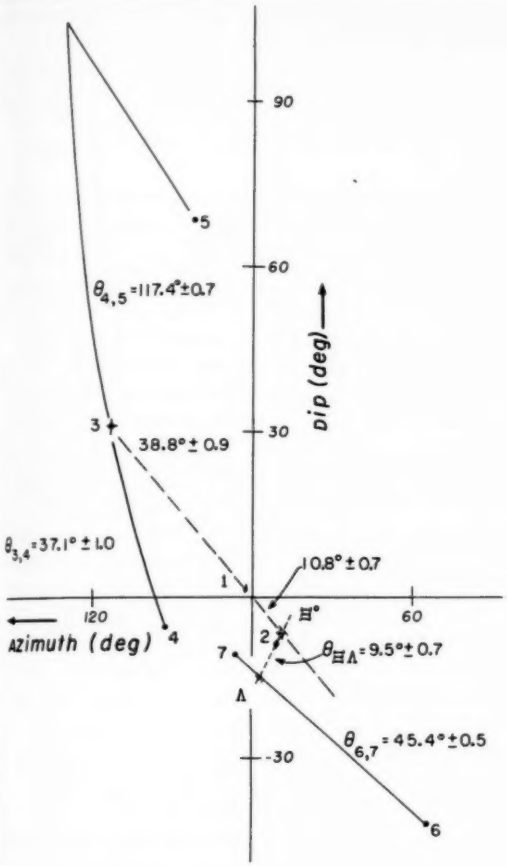


FIG. 2. Stereographic projection (Wulff plot) of the event. Observed tracks: 1, beam K^-; 2, line connecting end of beam track to vertex of Λ; 3, line connecting end of beam track to vertex of K^0; 4, π^-; 5, π^+; 6, π^-; 7, p. Inferred "tracks": Λ, obtained by balancing transverse momentum of Tracks 6 and 7; Ξ^0, obtained by intersection of production plane (containing Tracks 1 and 3) with the plane containing Track 2 and the Λ.

0.5 mm, requires stretching the errors by more than 5 standard deviations. Inelastic scatterings, for instance, $\Lambda+p\rightarrow\Sigma^0+p$, would also always give a visible recoil. Double scatterings, scatterings on deuterium, or neutron reactions on deuterium that might look like Λ events are exceedingly unlikely; for instance, the probability that we have in the entire experiment a double V^0 with a Λ scattering on deuterium that resembles a Ξ^0 event is $\leq 10^{-6}$.

The third possibility, a chance coincidence, can be shown to be most improbable on statistical grounds. Since the argument hinges on how well the event fits the production and decay of a Ξ^0 hyperon, let us now turn to this hypothesis. If we assume the K^0 meson to be produced in association with a heavy unstable particle, the incident particle being a beam K^- meson, then the extra energy available in the center of mass in the K^-+p system (compared with the π^-+p system) requires the heavy particle to be much heavier than a Σ^0. If this particle travels a distance of 3.7 cm and decays into a Λ and a π^0, then the presence of an associated Λ can be explained, as well as its apparent noncoplanarity.

The mass of the particle may then be deduced in two ways, i.e., from its production and from its decay. First, if we take the production process to be a two-body reaction, the heavy particle must lie in the plane formed by the beam track and the line of flight of the K^0 meson. Further, the direction of the heavy particle in this plane is fixed by the requirement that its path intersect that of the Λ. Then, using the production angles and the measured K^0 momentum, we can calculate the heavy-particle mass as well as the momentum of the incident K^- meson. The calculated momentum of the incident K^- is 1.13 ±0.06 Bev/c, which agrees well with the nominal beam momentum. (This serves as a first check on our hypothesis.) The heavy-particle mass is 1303 ± 28 Mev.

Second, if the heavy-particle velocity as determined in the above calculation is taken in conjunction with the observed Λ momentum and angle, a second mass determination is possible. The heavy-particle mass resulting from this calculation, based on the assumption of decay into a π^0 and a Λ, is 1349 ± 30 Mev. This value is insensitive to the heavy-particle velocity, and therefore the two determinations are nearly independent.

Combining the two mass determinations,[8] we obtain

$$M = 1326 \pm 20 \text{ Mev}.$$

The closeness of this result to the accepted Ξ^- mass of 1321 ± 35 Mev[9] is remarkable.

One might put the arguments the other way and ask to what extent the agreement (within errors) with the Ξ^- mass restricts the position, momentum, and angle of the decay Λ. In order for the Ξ^0 mass, as determined by its production, to vary by 30 Mev, the Λ need be moved (transversely) only 0.4 mm. Similarly, in order for the Ξ^0 mass, as determined by its decay, to

vary by 30 Mev, the Λ momentum must be changed by 50 Mev/c (at fixed angle), or the angle by 2° (at fixed momentum).

These restrictions form a strong argument against the possibility of accidental coincidence. A careful estimate shows that the probability of getting one such accidental event in the entire experiment is of the order of 10^{-5}. We have not been able to think of any more likely possibilities. Therefore, we believe that this event represents the production and decay of a Ξ^0, i.e. a hyperon of strangeness-2 and mass comparable to that of the Ξ^-.[10]

The measured dynamical variables of the event are:

K^0: Momentum, 277.5 ± 5.0 Mev/c; production angle (laboratory system), $38.8 \pm 0.9°$.

Λ: Momentum, 920 ± 50 Mev/c; angle (laboratory system) between Λ and Ξ^0, $9.5 \pm 0.7°$.

Ξ^0: Production angle (laboratory system), $10.8 \pm 0.7°$.

The incident K^- momentum agrees so well with the independently determined beam momentum, 1.15 ± 0.02 Bev/c, that it is highly probable that our K^- is one of the beam K^-'s. On this basis we can determine the mass much more precisely: $M_{\Xi^0} = 1308 \pm 8$ Mev at production. This gives $\chi^2 = 0.077$ ($\langle\chi^2\rangle = 1$).

If we consider all the information given by the production, the beam momentum, and the decay (assuming $\Xi^0 \to \pi^0 + \Lambda$), we find for the most probable mass $M_{\Xi^0} = 1311 \pm 8$ Mev. For this we find $\chi^2 = 1.45$ ($\langle\chi^2\rangle = 2$). The event cannot be used for a check of the decay mode; for instance, if we assume $\Xi^0 \to \gamma + \Lambda$, we find an even better fit ($\chi^2 = 0.247$).

The cross section, based on this one event, is $\sigma_{\Xi^0 K^0} \approx 50$ μb. We have not seen any examples of $K^- + p \to \Xi^- + K^+$; this sets a diffuse upper limit, $\sigma_{\Xi^- K^+} \leq 17$ μb. (No correction for lifetime is made here. If the lifetime of either Ξ is long compared with 5×10^{-10} sec, many would escape from the chamber.) Our one Ξ^0 lived 1.5×10^{-10} sec.

It is interesting to compare the above cross sections with those for the similar reactions

$$\pi^- + p \to \Sigma^- + K^+,$$

$$\pi^- + p \to \Sigma^0 + K^0,$$

at the same outgoing c.m. momentum (190 Mev/c)[11]:

$$\sigma_{\Sigma^- K^+} \approx 200 \ \mu\text{b},$$

$$\sigma_{\Sigma^0 K^0} \approx 400 \ \mu\text{b}.$$

At present the search for production of cascade hyperons in the 1.15-Bev/c K^- beam is being continued in collaboration with the Lawrence Radiation Laboratory 30-inch propane bubble chamber group.

The authors are greatly indebted to Dr. Edward J. Lofgren and the Bevatron staff for their help and cooperation, to J. Donald Gow and the bubble chamber crew for their successful operation of the bubble chamber, and to Glenn Eckman, Bob Watt, Bob Horne, and Ivan Muzinich for help in setting up the spectrometers. Dr. Bruce Cork, Dr. William Wenzel, and Dr. Glenn Lambertson gave valuable advice and encouragement with the spectrometers. Our scanning and measuring staff are to be thanked for their skilled work. George Edwards and William Salsig of the Mechanical Engineering group gave freely of their time in the spectrometer design and construction. Special thanks are due to Larry Ratner for his generous assistance with the magnetic measurements and to Dr. Frank Solmitz for help with the analysis.

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‡Presently at University of California at Los Angeles, Los Angeles, California.

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[10]A picture of a possible Ξ^0 was submitted by the Pic du Midi group to the Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958 (United Nations, Geneva, 1958).

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SEARCH FOR ENHANCEMENT OF BREMSSTRAHLUNG PRODUCED BY 575-Mev ELECTRONS IN A SINGLE CRYSTAL OF SILICON*

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Dyson and Überall[1] and Purcell[2] have predicted that considerable deviation from the Bethe-Heitler formulas for high-energy bremsstrahlung and electron pair production is expected in case the target material is crystalline. The basis of this effect is the low momentum transfer $\vec{q}$ to the Coulomb field in these processes. Bremsstrahlung by a 600-Mev electron leading to a 300-Mev photon leads to a momentum transfer to the Coulomb field down to a minimum value of 500 ev/c. Hence if such a transfer $\vec{q}$ corresponds to a reciprocal lattice vector of the crystal, the screened Coulomb field of the various atoms acts coherently and the production amplitude is enhanced. Detailed calculations of this effect have been carried out by Überall[3]; this experiment was aimed at checking these predictions.

A single-crystal plate of silicon of 0.013-in. thickness was placed in the analyzed beam of the Stanford electron linear accelerator. The crystal was mounted in a double goniometer to permit rotation of the crystal about two coplanar perpendicular axes perpendicular to the beam. The zero position of the goniometer was adjusted to correspond to the (100) plane of the crystal perpendicular to the beam; this orientation was established by Laue back-reflection photographs using a standard metallographic camera.

X-rays radiated in the crystal were detected by letting them produce photopions of specified energy in a polyethylene target. The positive pions were counted in a scintillation counter after magnetic analysis via the μ-decay positrons. The counting rate was examined as a function of "scanning" the goniometer over a range of approximately ± 0.04 radian about both axes in steps of approximately 0.010 radian. About 1500 counts per point were taken. No statistically significant dependence of x-ray intensity on crystal orientation was found.

This result was assessed quantitatively by folding the functions computed by Überall into the angular resolution of the experiment as defined by the multiple scattering of the primary beam in the crystal target itself. Since this scattering involves small momentum transfers also, one could expect crystal coherence effects here as well, and therefore measured rather than computed scattering angles were used in the computation of the "expected" x-ray intensity variation as a function of angle between the crystal planes and the incident electron beam angle.

Figure 1 is a map of the counts obtained as a

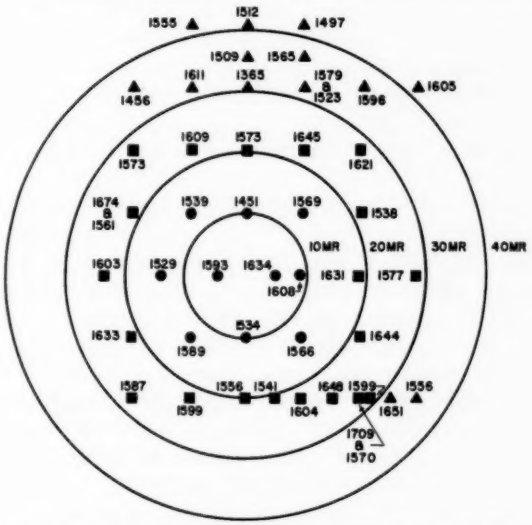

FIG. 1. The circles shown are lines of equal angle between the direction of the incident electron beam and the [100] axis of the Si crystal in which bremsstrahlung is produced. The angles are shown in milliradians (mr). Each point represents the measured intensity in counts of bremsstrahlung near an energy of 235 Mev.

function of crystal orientation about two orthogonal axes. Drawn on this figure also are circles of expected equal intensity. Clearly no significant correspondence exists. Qualitatively, the mean value of the central points (circles), in a zone of 8×10^{-3} radian mean angular displacement, is 1556 ± 13; the average of the points (squares) in a zone of mean angular displacement of 0.02 radian, is 1602 ± 9; and the average of points (triangles) in a zone of mean angular displacement of 0.035 radian is 1541 ± 10. The "expected" counts should be in the ratio 22: 19.5: 17.4, using the Überall calculation for a crystal temperature at absolute zero.

We find it difficult to account for this discrepancy. The effect of lattice vibration has been discussed by Überall, and on the basis of his calculations is insufficient to suppress the effect to the extent needed to fit the data. More detailed calculations on this point are in progress.[4] The question whether sufficient radiation dislocation could have occurred during bombardment has been examined; assuming an activation energy of 25 ev for a dislocation, only 5×10^{-6} of the Si atoms would have been affected. The crystals were furnished to us via Professor J. W. M. DuMond, California Institute of Technology, by Mr. W. R. Runyon of the Texas Instrument Company, Dallas. Mr. Ronald Willems, California Institute of Technology, kindly ran x-ray "rocking curves" on the crystals used: the angular half-width at half-maximum is only 8×10^{-4} radian; this is negligible for our considerations. Laue back-reflection pictures taken before and after our bombardments showed no observable deterioration.

We are unable to suggest an explanation for this result other than possibly the lack of validity of some of the approximations made in the calculations on the effect of lattice vibrations, or possibly of the Born approximation.

We are greatly indebted to Mr. Richard Bush, of the Stanford Department of Metallurgical Engineering, for the use of the Laue camera. We are also indebted to Professor J. W. M. DuMond for furnishing the crystals and for valuable advice. We wish to thank Dr. G. L. Pearson of Bell Telephone Laboratories, for informing us about the techniques of lapping single crystals of Si and their chemical etching.

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[4]L. I. Schiff (private communication).

NONMESONIC/MESONIC DECAY RATIO OF HELIUM HYPERFRAGMENTS*

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In order to understand the decay interactions of the bound Λ^0,[1-6] it is important that the nonmesonic/mesonic ratio (Q) in the decay of hyperfragments be experimentally well determined. Previously reported experimental values of Q for ${}_\Lambda$He (2.3 ± 1.0; 1.1 ± 0.5)[7,8] disagree with the existing calculations[4,5] ($4\leq Q\leq10$). Notwithstanding the present uncertainties in the theoretical work, which have been recently discussed by Dalitz,[5] certain refinements in the experimental work are desirable. It is known that K^- capture stars are prolific sources of hyperfragments. Thus problems of contamination in the event samples are diminished and the increased number of events allow the results of kinematic analysis of the nonmesonic events to be studied. Track thickness ionization measurements in the Ilford fine-grain $K5$ and $L4$ emulsion allow the use of a rigid event-selection criterion.

In this emulsion experiment we have obtained 33 examples of nonmesonic decay of ${}_\Lambda$He, which were produced in K^--capture stars, had ranges $\geq59\ \mu$, and decayed with the formation of two visible prongs. (The K^- mesons were produced at the Bevatron; they interacted at rest.) The 33 events were selected by means of track-thickness[9] (profile) and gap-count measurements on the connecting tracks. The results of the profile calibration measurements made on known charge 1 and 2 tracks are shown in Fig. 1. The dotted line is seen to separate the $Z=1$ and $Z=2$ calibration points at all dip angles to the extent that an individual $Z=1$, 2 track can be charge identified to ~95% level-of-confidence. No $Z=3$ calibration was necessary since the range distribution of the

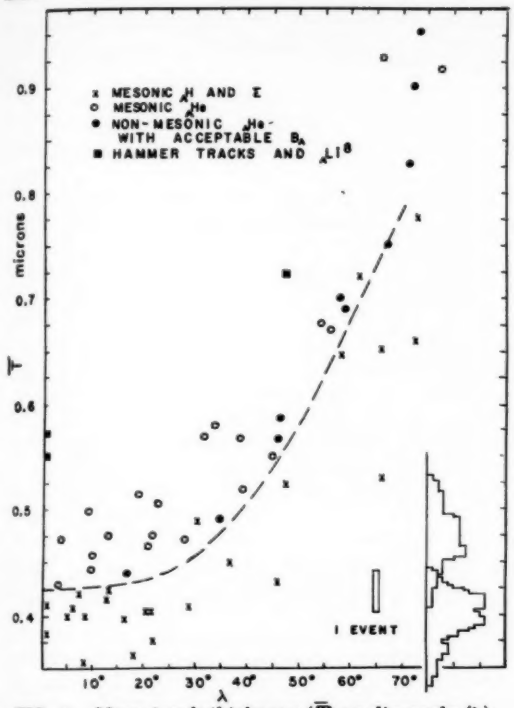

FIG. 1. Mean track thickness ($\overline{T}$) vs dip angle (λ) for known charge 1, 2 tracks in Ilford $\overline{K}5$ and $L4$ emulsion. Thickness measurements were made at $\sim 5000\times$ magnification with 0.5 μ projected cell lengths from 25 μ to (37/cosλ) μ residual range. The dashed line shows the optimum separation between charge 1 and 2 tracks. No dependence of $\overline{T}$ on depth in emulsion was found. The insert shows the $\overline{T}$ distributions for charge 1 and 2 tracks with $\lambda < 30°$.

connecting tracks of mesonic ${}_\Lambda$Li events from K^--capture stars at rest is known[10] to be < 50 μ.

In the emulsion scannings in which these hyperfragments were found, no systematic search for possible one-prong nonmesonic decays of ${}_\Lambda$He was carried out. Nevertheless, one probable example of ${}_\Lambda He^5 \rightarrow He^4 + n$ has been found in this stack. The identification of this event is based on profile measurements of the stopping connecting track (range = 128 μ, dip angle = 6°) and of the single prong, which has the characteristic 495-μ length required for the proper energy release (dip angle of this prong is 3°).

The 33 two-prong nonmesonic events are to be compared with 22 examples of π^--mesonic decay of ${}_\Lambda He^{4,5}$ with connecting tracks ≥ 50 μ found in this stack. These mesonic decays have already been reported as part of the EFINS-NU (Enrico Fermi Institute of Nuclear Studies—Northwestern University) collaboration experiment[10] on π^--mesonic hyperfragments. It should be pointed out that on the basis of the large energy releases, none of the 33 nonmesonic events could be interpreted as π^0-mesonic hyperfragment decays. A scanning bias check shows the over-all efficiencies for finding the two-prong nonmesonic and π^--mesonic events in the experiment to be $\epsilon_{\text{nonmes.}} \tilde{=} \epsilon_{\text{mes.}} \tilde{=} 80\%$. Thus the experimental ratio determined here is $Q' = 33/22 = 1.5 \pm 0.4$. Since the one-prong nonmesonic decays were not included in this experiment, this value must be considered as a lower limit to the total nonmesonic ratio for ${}_\Lambda$He.

An attempt was made to estimate the relative frequencies of the various nonmesonic decay modes of ${}_\Lambda He^{4,5}$ in our sample. With the use of an IBM-650 hyperfragment analysis program,[11] the two-prong nonmesonic events considered in this experiment were subjected to the following analysis.

For each event, the residual momentum of the two charged particles (all permutations of $Z = 1$, $A = 1, 2, 3$ prong identifications were assumed) was attributed to one neutron and the binding energy of the Λ^0 (B_Λ) was then calculated for that assumption. For each event the resultant B_Λ closest to the known mean B_Λ from the study of mesonic ${}_\Lambda$He decays[10] (2.3 Mev) was selected. The dotted curve in Fig. 2 is the expected B_Λ distribution based on the mean B_Λ for ${}_\Lambda He^{4,5}$ and a standard deviation of 3.2 Mev resulting from the experimental range straggling and angle-measurement uncertainties in the nonmesonic events. 23 events have B_Λ's which fall within the 98% area limits of the Gaussian. Those events are considered to be kinematically identified, and their B_Λ's are plotted in Fig. 2. All the possible nonmesonic decay modes of ${}_\Lambda He^{4,5}$ are shown in Table I with the number of identified events of each mode in our sample. From the B_Λ distribution of the 10 "background" events (i.e., those with $B_\Lambda < -5.6$ or $B_\Lambda > 10.4$ Mev), it can be estimated that 3 of the 23 identified events may be misidentified. The 10 "background" events are assumed to have 2 or 3 neutrons emitted and thus are examples of modes $c5$, $c5'$, or $c4$.

In accordance with the recent work of Baldo-Ceolin et al.[12] and Ferrari and Fonda,[6] modes $a5$ and $a4$ in Table I are referred to as proton-stimulated decays and modes $e5'$ and $c4$ as neutron-stimulated decays [i.e., $\Lambda^0 + (p \text{ or } n) \rightarrow n + (p \text{ or } n)$]. The prong momentum distributions in the decay modes $a5$ and $a4$ (see Fig. 3) are of interest, for we see that the proton and neutron

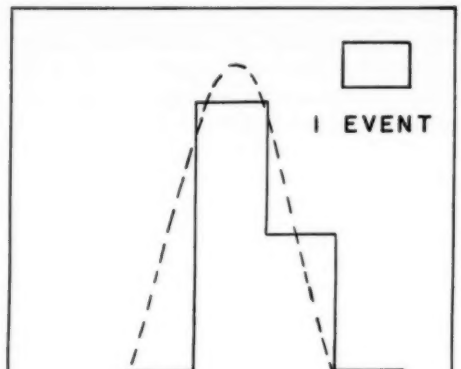

FIG. 2. B_Λ distribution for those ${}_\Lambda$He events which furnish acceptable binding energies (see text). The dashed curve is the expected distribution based on the known B_Λ for ${}_\Lambda He^{4,5}$ and a 3.2-Mev standard deviation caused by range straggling and angle-measurement uncertainties in the nonmesonic events.

momentum distributions are similar. This is consistent both with our identification of these events and with the single nucleon stimulation picture.

Those events which decay via mode $b5$ do not represent simple cases of stimulated decay of the Λ^0. They can be interpreted either in terms of a two-nucleon stimulation process or of the existence of the intermediate states H^3+H^1+n or He^3+2n with subsequent final-state interaction[13] producing $2H^2+n$.

Table I. Nonmesonic decay modes of ${}_\Lambda He^{4,5}$. The numbers in parentheses refer to events for which no single acceptable B_Λ results from the kinematic analysis. Except for the one indicated event, modes $e4$, $e5$, $e5'$ were excluded from consideration in this experiment.

	${}_\Lambda He^5$	Number of events			${}_\Lambda He^4$	
($a5$)	H^3+H^1+n	8	(2)	8	H^2+H^1+n	($a4$)
($b5$)	H^2+H^2+n	5				
($c5$)	H^2+H^1+2n		(10)		H^1+H^1+2n	($c4$)
($c5'$)	H^1+H^1+3n					
($d5$)	H^3+H^2	0		0	H^3+H^1	($d4$)
				0	H^2+H^2	($d4'$)
($e5$)	He^4+n	1		...	He^3+n	($e4$)
($e5'$)	He^3+n+n					

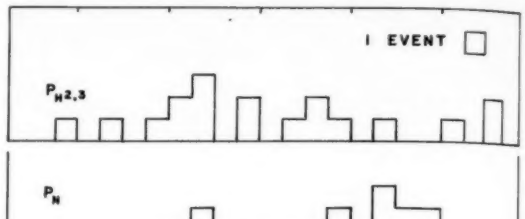

FIG. 3. The particle momentum distributions in those events which were kinematically identified as ${}_\Lambda He^{4,5} \rightarrow H^{2,3}+H^1+n$.

The inclusion of the one-prong nonmesonic decays of ${}_\Lambda He^{4,5}$ in an experimental determination of the total nonmesonic decay rate is a necessary consideration in the planning of future hyperfragment experiments. In particular, mode $e5'$ must be included in an experimentally determined neutron/proton stimulation ratio. The identification of this neutron-stimulated decay mode of ${}_\Lambda He^5$ presents a difficult problem, however, for the He^3 prong can have a wide range of values. If the recoil (He^3) momentum distribution in these decays is similar to the recoil ($H^{2,3}$) momentum distribution in the proton-stimulated decays of ${}_\Lambda He^{4,5}$, then the $P_{H^{2,3}}$ distribution in Fig. 3 tells us that ~80% of the He^3 prongs in mode $e5'$ will have ranges $\geq 70\,\mu$. The appearances of modes $e4$ and $e5$ are self-evident because of the characteristic prong lengths and should provide no particular difficulties. The one example of $e5$ reported here and a possible example of $e4$ reported by Silverstein[8] point out the existence of these modes.

I am grateful to Professor J. Roberts, Professor L. Brown, and Professor M. Peshkin for consultation and continued encouragement, and to Dr. S. Limentani and Phillip Steinberg for collaboration on the mesonic decay work and in setting up the profile-measuring microscope.

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NATURE OF THE VECTOR INTERACTION IN μ^--ABSORPTION*

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Feynman and Gell-Mann have suggested[1] that the strangeness-conserving vector lepton interaction current $J_\lambda^{(V)}$ is conserved, and equal (up to a constant C_V/e, and an isospin rotation) to the isovector part of the electromagnetic current. The β-decay experiments that have been suggested that might test this idea are either very difficult because the "weak-magnetism" effects are very small and easily masked by Coulomb and other effects,[2,3] or do not unambiguously distinguish between the new theory and any reasonable model of nuclear decay.[4] Therefore, one is naturally led to consider, instead of β-decay, the very high momentum-transfer process of μ^- absorption.[5] We wish to suggest that a comparison of the rate of any $0\to0$ *no* μ^- capture, with the cross section for the corresponding inelastic electron scattering process, may serve as a definitive test of the Feynman–Gell-Mann proposal.

In general, the transition probability for a process $\mu^- + A \to \nu + B$ will depend on matrix elements of the vector and the axial-vector currents. In order to test a theory of the vector interaction it is necessary to consider the case where A, B have zero spin and equal parity so that the axial current cannot contribute. Since A must be fairly stable it must necessarily be the ground state of an even-even nucleus; then B is some state of an odd-odd nucleus. The total rate for the particular transition $A\to B$ is given by

$$\omega = \frac{2Z_A{}^3\alpha^3 m_\mu{}^3 m_A{}^4 P_\nu{}^4}{(2\pi)^6(m_\mu+m_A)^4(q^2)^2}\,|F(q^2)C_V\sqrt{2}|^2, \qquad (1)$$

where the function $F(q^2)$ is defined, for conserved currents, by

$$\langle B|J_\lambda^{(V)}|A\rangle = (2\pi)^{-3}(4m_A E_B)^{-1/2}F(q^2)C_V\sqrt{2}$$
$$\times\{(P_B+P_A)_\lambda - [(m_A{}^2-m_B{}^2)/q^2](P_B-P_A)_\lambda\}, \qquad (2)$$

and q^2 is the invariant momentum transfer,

$$q^2 \equiv (P_A-P_B)^2 = m_\mu{}^2 - \frac{m_\mu(m_\mu{}^2+m_B{}^2-m_A{}^2)}{m_\mu+m_A} \sim (100\ \mathrm{Mev}/c)^2. \qquad (3)$$

Now suppose that A is an isospin singlet, that B has $T=1$, $T_3=-1$, and that A^* is the excited state of A belonging to the same triplet as B. The inelastic electron scattering process, $e + A \to e + A^*$, has a matrix element given by the Feynman–Gell-Mann theory as

$$\langle A^*|J_\lambda^{(\mathrm{el})}|A\rangle = (2\pi)^{-3}(4m_A E_{A^*})^{-1/2}F(q^2)e$$
$$\times\{(P_{A^*}+P_A)_\lambda - [(m_A{}^2-m_{A^*}{}^2)/q^2](P_{A^*}-P_A)_\lambda\}, \qquad (4)$$

where $J_\lambda^{(\mathrm{el})}$ is the electric current, $q^2=(P_A-P_{A^*})^2$, and $F(q^2)$ is the same as in (2). The differential cross section is then given by

$$\left(\frac{d\sigma}{d\Omega}\right)_{\mathrm{lab}} = \frac{\alpha^2\cos^2(\theta/2)\,|F(q^2)|^2}{4E_e[E_e+(m_A{}^2-m_{A^*}{}^2)/2m_A]\sin^4(\theta/2)[1+(2E_e/m_A)\sin^2(\theta/2)]}, \qquad (5)$$

where we use the Born approximation, and neglect m_e. Clearly, if the electron scattering experiment is performed with energy and angle chosen so that q^2 is given by (3), the ratio of (5) and (1) becomes a known function independent of $F(q^2)$;

$$\omega/\left(\frac{d\sigma}{d\Omega}\right)_{\text{lab}} = \frac{Z^3 m_\mu{}^5 m_A{}^4 P_\nu{}^4 [1+(2E_e/m_A)\sin^2(\theta/2)]^{\frac{1}{2}} \ln 2/(ft)_{O^{14}}}{64\pi^3(137)m_e{}^5(m_\mu+m_A)^4 E_e[E_e+(m_A{}^2-m_{A^*}{}^2)/2m_A]\cos^2(\theta/2)}. \quad (6)$$

This prediction seems to be the only test of the Feynman–Gell-Mann theory for muon interactions that is completely free of theoretical ambiguities arising from our ignorance of the details of nuclear structure.

The really difficult part of this experiment will be the measurement of the $\mu^- + A \to \nu + B$ absorption rate. Light odd-odd nuclei are never found with zero-spin ground states so that B will have to be an excited state, which must be detected by observing the de-excitation γ-rays. (The only suitable nuclear level known to the author occurs in Al^{28}.) The absorption rate is likely to be rather low, since by isotopic spin conservation $F(0)=0$; if A, B were members of the same triplet we would have $F(0)=1$, but then the elastic electron scattering $A+e \to A+e$ would depend on the isoscalar, as well as the isovector, part of the electric current. It is of course desirable to perform this experiment in as light a nucleus as possible, to avoid possible departures from isospin conservation and from the Born approximation in electron scattering, and to have the nuclear energy levels well separated.

Similar remarks also apply in the case of nonzero spins, provided that a reliable theory of the axial-vector contributions to μ^- capture may be developed.

I wish to thank Professor H. Kendall, Professor T. D. Lee, and Professor H. Primakoff for valuable comments.

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EQUIVALENCE TRANSFORMATION OF THE β-DECAY OPERATOR

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There have been attempts recently[1,2] to use the Chew-Low fixed source theory with cutoff for calculating mesonic corrections to the β-decay and magnetic moment of the nucleon. This method uses pseudovector coupling in no-recoil approximation and with neglect of virtual pair creation. Therefore it must be interpreted as the nonrelativistic limit of the theory obtained after application of the equivalence transformation[3] to the nucleon and pion fields coupled by the pseudoscalar interaction. In particular, the nucleon field ψ transforms into ψ':

$$\psi \to \psi' = \exp(iS)\psi, \quad (1)$$

where

$$S = (\vec{\tau}\cdot\vec{\phi})\gamma_5 \tan^{-1}(G\phi/m)/2\phi. \quad (2)$$

In Eq. (2) we have defined

$$\phi \equiv [\vec{\phi}\cdot\vec{\phi}]^{1/2}. \quad (3)$$

Since the nucleons one deals with in Chew-Low theory are described not by ψ but by ψ', it follows that the operators coupling nucleons to other fields (lepton and photon) must also be transformed:

$$\psi^* Q \psi \to \psi^* e^{-iS} Q e^{iS} \psi. \quad (4)$$

We have been interested especially in the mesonic corrections to the axial vector β-decay operator.[4] Under the equivalence transformation, and in the nonrelativistic limit for the nucleons,

this transforms as follows:

$$Q = \gamma_4\gamma_5\vec{\gamma}\tau_x - Q',$$

$$Q' = Q\left\{1 + \frac{[\vec{\phi}\times(\vec{\phi}\times\vec{\tau})]_x}{\phi^2\tau_x}\left[1-(1+G^2\phi^2/m^2)^{-1/2}\right]\right\}. \quad (5)$$

The operator $\delta Q = Q' - Q$ in (5) cannot meaningfully be written as a unique power series in G. However, we have developed a method for determining exactly its matrix elements between any given approximation to the physical nucleon state vectors. This exact treatment of δQ is further motivated by the observation that the effects we are considering involve diagrams where the mesons are emitted and re-absorbed at one and the same point on the nucleon propagator (simultaneously with β-emission). Thus Chew-Low theory provides no reason to believe that a perturbation treatment of this particular series of diagrams will converge at all fast, since they do not depend on the meson-admixture coefficients. In fact we know a posteriori that although the effect is finite in the no cutoff limit, each separate term of the series then diverges. Thus we are allowing for a slowly convergent set of contributions exactly but must still deal with the more quickly convergent Chew-Low perturbation series term by term.

To illustrate the procedure we sketch the calculation for the bare nucleon state, which amounts to evaluating the meson-vacuum expectation value of δQ. Consider the first (G-independent) term of δQ:

$$\frac{[\vec{\phi}\times(\vec{\phi}\times\vec{\tau})]_x}{\phi^2\tau_x} = \frac{[(\vec{\phi}\cdot\vec{\tau})\phi_x - (\vec{\phi}\cdot\vec{\phi})\tau_x]}{\phi^2\tau_x}.$$

It clearly has vacuum expectation value $-2/3$, since only those terms of the numerator contribute which are even in the three components of $\vec{\phi}$ taken separately.

Next, in the sense of generalized function theory,[5] we decompose the remaining operator Δ into a three-dimensional Fourier integral:

$$\Delta = -Q\int f_x(\vec{\alpha})\exp(i\vec{\alpha}\cdot\vec{\phi})d^3\alpha, \quad (6)$$

where

$$f_x(\vec{\alpha}) = (2\pi)^{-3}\int\frac{[\vec{\phi}\times(\vec{\phi}\times\vec{\tau})]_x}{\phi^2\tau_x}(1+G^2\phi^2/m^2)^{-1/2} \times \exp(-i\vec{\alpha}\cdot\vec{\phi})d^3\phi. \quad (7)$$

$\vec{\phi}$ is now written in terms of its plane wave components with a momentum cutoff K. By exploiting the analogy with the simple harmonic oscillator in ordinary quantum mechanics, one finds the vacuum expectation value

$$\langle 0|\exp(i\vec{\alpha}\cdot\vec{\phi})|0\rangle = \exp(-\alpha^2 C/2), \quad (8)$$

where

$$C = V^{-1}\textstyle\sum_{k<K}(1/2\omega_k). \quad (9)$$

The sum in the exponent on the right-hand side of (8) arises because the vacuum state vector is a product of oscillator ground states for each of the field oscillators. Note that $C(K)$ is essentially the commutator of the annihilation and creation operator parts of $\vec{\phi}$.

Inserting (8) into the vacuum expectation value of (6), and using the definition (7), we get

$$\Delta = -Q\int d^3\alpha\exp(-\alpha^2C/2)\times(2\pi)^{-3}\int\frac{[\vec{\phi}\times(\vec{\phi}\times\vec{\tau})]_x}{\phi^2\tau_x}(1+G^2\phi^2/m^2)^{-1/2}\exp(-i\alpha\cdot\phi)d^3\phi. \quad (10)$$

But in the limit $K\to\infty$, $C\to\infty$, and the integral above vanishes, giving the final result

$$\langle 0|Q'|0\rangle = \langle 0|Q(1-\tfrac{2}{3})|0\rangle = \langle 0|Q|0\rangle/3.$$

We have made a similar calculation allowing one-meson admixtures to the bare nucleon states, and again found that the contribution of δQ is $-2/3$ that of Q. We have not succeeded so far in generalizing this reduction factor to arbitrary nucleon states.

For finite K in the vicinity of the current value, the reduction factors are not significantly affected.

Similar results are obtained for the nucleon contribution to the isotopic vector part of the magnetic moment.

We conclude that the Chew-Low method, when applied consistently, cannot adequately deal with effects to which the nucleon core itself contributes significantly. In other words, though the method may yield a good description of the tail

of the physical nucleon, it fails near the core where nuclear β-decay and Dirac moment are concentrated.

A more detailed account of this work will be published elsewhere.

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EVIDENCE FOR RESONANCE STATES IN THE K^--N SYSTEM

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(Received January 19, 1959)

One of the striking features of σ_{el} for K^--p scattering is its consistently high value compared with σ_{el} for K^+-p scattering,[1] [$\sigma_{el}(K^-p)/4\pi\lambda\bar{}^2 \sim \frac{1}{8}$, $\sigma_{el}(K^+p)/4\pi\lambda\bar{}^2 \sim \frac{1}{34}$ for T_K(lab) ≈ 25 Mev], near threshold. This, combined with the suspicion that the K^-p "potential" is attractive and the K^+p "potential" repulsive at low energies, and a characteristic peaking[2] of $\sigma_{el}(K^-p)$ around T_K(lab) $= 25$ Mev, seems to suggest the possibility that interactions in the energy range 15-60 Mev proceed mainly through a $J=\frac{1}{2}$ resonance centered around $T_K = 25$ Mev. Such a resonance state (to be called Σ^*) would be expected to show itself both in elastic scattering as well as in the K^--absorption processes. Assuming the existence of a resonance at E_0 (c.m. energy), we use the one-level resonance formulas

$$\sigma_{el} = \frac{(2J+1)}{2} f_e(I) \frac{\pi}{k_K^2} \frac{\Gamma_K^2}{(E-E_0)^2 + \frac{1}{4}(\Gamma_K+\Gamma_\pi)^2}, \quad (1)$$

$$\sigma_{abs}(\pi^\pm, \Sigma^\mp) = \frac{2J+1}{2} f_a(I) \frac{\pi}{k_K^2} \frac{\Gamma_K \Gamma_\pi}{(E-E_0)^2 + \frac{1}{4}(\Gamma_K+\Gamma_\pi)^2}. \quad (2)$$

Here J is the angular-momentum of the resonance state and $f(I)$ are the isotopic spin factors:

$$f_e(I) = \tfrac{1}{4} \text{ for } I=0, 1;$$

$$f_a(I) = \tfrac{1}{3} \text{ for } I=0$$
$$= \tfrac{1}{2} \text{ for } I=1;$$

and Γ_K and Γ_π are the partial widths. From (1) and (2),

$$\sigma_{el}/\sigma_{abs} = f_e(I)\Gamma_K / f_a(I)\Gamma_\pi. \quad (3)$$

The energy dependence of Γ_K and Γ_π is sensitive to the $(K\Sigma)$ and $(\Sigma^*\Sigma)$ parities. If we assume that both the $(\Sigma\pi)$ and (Kp) systems are in S states, with momentum k in the center-of-mass system, we may take

$$\Gamma = 2ky^2, \quad (4)$$

where y is an energy-independent parameter. This implies (with the convention that N and Σ have the same parity "plus") that both Σ^* and K have intrinsic parity "minus." Relation (3) then reduces to

$$\frac{k_K \sigma_{abs}}{k_\pi \sigma_{el}} = \text{const} = \frac{f_a(I) y_\pi^2}{f_e(I) y_K^2}. \quad (5)$$

This checks fairly well with the experimental results in the energy region $15 < T_K < 60$ Mev, provided we use the Berkeley bubble-chamber data (rather than the emulsion data) for $25 < T_K < 35$ Mev.

If we have P waves, then

$$\Gamma = \frac{2ky^2}{1+(ka)^{-2}} \approx 2k^3 a^2 y^2, \ (ka \ll 1), \quad (6)$$

where a is the interaction radius. The possibility that the reactions go purely through P waves is excluded by the relation corresponding to (5), and a less satisfactory fit is obtained with the assumption that just the $\pi\Sigma$ system is in a P state.

Noting that

$$\sigma_{el}(E=E_0) = \frac{k^2}{4\pi} f_e(I) \left[1 + \frac{k_\pi y_\pi^2}{k_K y_K^2}\right]^{-1}_{E=E_0}, \quad (7)$$

and combining (6) and (7) one should also be able to deduce the isotropic spin value of Σ^*. Unfortunately the data are not good enough to distinguish the case $I=0$ from $I=1$. We shall come back to this later.

However, with the spin and parity assignment $(\frac{1}{2}, -)$ to the Σ^* state, we can now correlate the following facts:

(1) The angular distribution for

$$K^- + p \rightarrow \Sigma^* \rightarrow K^- + p$$

must be isotropic. This result depends on the

fact that $J=\frac{1}{2}$.

(2) The angular distribution for

$$K^-+p\rightarrow\Sigma^*\rightarrow\pi+\Sigma$$

is isotropic. This is also because $J=\frac{1}{2}$ and does not depend on the parity assignment of Σ^*. Both (1) and (2) seem borne out experimentally.

(3) The K^-p "potential" must be attractive. In a sense this is what first led to a consideration of the resonance hypothesis. Given the resonance hypothesis one can formally check back (using the usual computational approximations) that the effective "potential" is indeed attractive. So far as we know, this is the only "explanation" of this experimentally suspected result. Without the resonance hypothesis, the K^-p "potential" turns out to be repulsive both for scalar and pseudoscalar K-mesons in a lowest order calculation.

(4) In the production processes

$$\begin{aligned}\pi^-+p&\rightarrow K^++\Sigma^-, &&\text{(i)}\\ &\rightarrow K^0+\Sigma^0, &&\text{(ii)}\\ &\rightarrow K^0+\Lambda^0, &&\text{(iii)}\\ \pi^++p&\rightarrow K^++\Sigma^+, &&\text{(iv)}\end{aligned}$$

it is easy to verify that the existence of Σ^* effects only (i) and (iv), in the same approximation used in (3) (treating Σ^* as a particle and working with lowest order graphs[3]). It is a long-standing problem that hyperons seem to be emitted in the backward direction in the c.m. system in processes (i) and (iv) and in the forward direction in (ii) and (iii). The existence of Σ^* provides a natural mechanism for differentiating (i) and (iv) from (ii) and (iii).

Next we turn to the question of the isotopic spin of Σ^*. For pure $I=0$ or $I=1$ states the cross sections should be in the ratios indicated in Table I. Experimentally, for $20<T_K<60$ Mev, the production ratio on protons is

$$\Sigma^+/\Sigma^-\approx 1,$$

while nothing is known about neutral hyperon production. Also

$$\sigma(K^-+p\rightarrow n+K^0)/\sigma(K^-+p\rightarrow p+K^-)>\tfrac{1}{5}.$$

This is not sufficient to fix the I value. A definite assignment would be possible, if one could determine whether Σ^0 or Λ^0 are produced predominantly in K^-p absorption in this energy region, indicating $I=0$ or 1, respectively.

Table I. Predicted cross-section ratios for pure $I=0$ or $I=1$ states.

	$I=0$	$I=1$
$K^-+p\rightarrow p+K^-$	a	α
$\rightarrow n+K^0$	a	α
$\rightarrow\Sigma^++\pi^-$	b	β
$\rightarrow\Sigma^-+\pi^+$	b	β
$\rightarrow\Sigma^0+\pi^0$	b	0
$\rightarrow\Lambda^0+\pi^0$	0	γ
$K^-+n\rightarrow n+K^-$	0	4α
$\rightarrow\Sigma^-+\pi^0$	0	2β
$\rightarrow\Sigma^0+\pi^-$	0	2β
$\rightarrow\Lambda^0+\pi^-$	0	2γ

So far we have only considered the Σ^-/Σ^+ ratio in the region $T_K>20$ Mev. We now consider it near threshold. According to the one-level picture presented above, this ratio should be unity. Experimentally, for $T_K=0$ this is nearer two; for increasing energy ($T_K\sim 10$ Mev) it rises to seven and then around 20 Mev it settles down to ~1. This behavior seems to call for a very special type of interference with amplitudes of opposite isotopic spin assignment in the energy range $0<T_K<20$ Mev. We have considered various simple models for this other amplitude which could give the experimental behavior in this region, while the resonance dominates for $T_K>20$ Mev. Considerations on this will be published in a separate note with J. Tiomno.

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[1]All data in this note are taken from the report given by M. F. Kaplon, at the 1958 Annual International Conference on High-Energy Physics at CERN, edited by B. Ferretti (Cern, Geneva, 1958), p. 171.

[2]See Fig. 8 and Fig. 13 of reference 1. Note that the simpler interpretation in terms of complex scattering lengths [see R. H. Dalitz, 1958 Annual International Conference on High-Energy Physics at CERN, edited by B. Ferretti (CERN, Geneva, 1958), p. 187] predicts a steady rising of the cross section as one approaches threshold. Both bubble-chamber and emulsion data suggest a fall, though the errors are large.

[3]The assumed interactions $\Sigma^*\rightarrow p+K^-$, $\Sigma^*\rightarrow\Sigma^\pm+\pi^\mp$ imply $\Sigma^*+K^+\rightarrow p$, $\Sigma^*+\pi^\pm\rightarrow\Sigma^\pm$.

ABSTRACTS

In this section are printed the abstracts of Articles that have been forwarded to The American Institute of Physics for publication in THE PHYSICAL REVIEW. In quoting information obtained from this section before the appearance of the corresponding Article, reference should be made to "Physical Review (to be published)" rather than to this Journal.

LOW TEMPERATURE BEHAVIOR OF A DILUTE BOSE SYSTEM OF HARD SPHERES. II. NON-EQUILIBRIUM PROPERTIES. T. D. Lee, Columbia University, New York, New York, and C. N. Yang, Institute for Advanced Study, Princeton, New Jersey (Received November 3, 1958).

The calculation of a previous paper is extended to cover nonequilibrium properties. The phenomena of superfluidity, critical velocity, and "infinite heat conductivity" are given natural explanations. By using classical kinetic theory on the wave packets, hydrodynamical equations for reversible flow is derived and the dependence of the two "sound velocities" on the temperature studied. The relationship between macroscopic sound vibrations and microscopic excitations is analyzed. The work is confined to the model of a dilute hard-sphere Bose system.

ELECTRON MOBILITY IN PARTIALLY IONIZED ATOMIC HYDROGEN. Mahendra Singh Sodha, Physics Research Department, Armour Research Foundation, Chicago, Illinois (Received September 17, 1958; revised manuscript received November 7, 1958).

Hall and drift mobilities of electrons in partially ionized atomic hydrogen have been investigated, taking into account the scattering by ions, neutral atoms, and other electrons.

LANGEVIN EQUATION AND THE AC CONDUCTIVITY OF NON-MAXWELLIAN PLASMAS. P. Molmud, Space Technology Laboratories, Los Angeles, California (Received November 3, 1958).

The use of the Langevin equation, $d\vec{v}/dt + g\vec{v} = (q/m)[\vec{E}_0 e^{i\omega t} + (1/c)\vec{v}\times\vec{H}]$, to describe the electrical conductivity of a non-Maxwellian plasma (a weakly ionized gas in which the average electron collision frequency is temperature dependent) may be in error unless it is understood that the dissipative term, g, is complex. In the limiting cases of either high or low pressures the imaginary part of g is negligible. The real and imaginary parts of g are evaluated for these limiting cases, for four different gases: air, helium, Maxwellian gas, and water. The real part of g is shown to be the average collision frequency multiplied by a numerical factor, the size of which depends on the nature of the gas and the pressure limit.

VELOCITY OF SOUND IN LIQUID He^3. K. R. Atkins and H. Flicker, Department of Physics, University of Pennsylvania, Philadelphia, Pennsylvania (Received October 20, 1958).

The velocity of sound in liquid He^3 along the vapor pressure curve between 1.2°K and 3.2°K has been measured using a pulse technique at 14 Mc/sec. Between 2°K and 2.6°K the variation with pressure up to 1 atmosphere has been investigated. The results are combined with other known data to calculate various thermodynamic quantities for the liquid. Available data on the specific heat along the vapor pressure curve have been converted into the more theoretically interesting case of the specific heat at constant volume along an isopycnal. The phonon contribution to the specific heat is discussed.

EXCITATIONS IN LIQUID HELIUM: NEUTRON SCATTERING MEASUREMENTS. J. L. Yarnell, G. P. Arnold, P. J. Bendt, and E. C. Kerr, Los Alamos Scientific Laboratory of the University of California, Los Alamos, New Mexico (Received October 16, 1958).

The energy-momentum spectrum of the excitations in liquid helium II has been measured in a neutron scattering experiment. At $T = 1.1$°K, data were obtained in the momentum range $p/\hbar = 0.55$ to $2.36\,A^{-1}$. The results bear a striking resemblance to the spectrum proposed by Landau in 1947. If a phonon spectrum of the form $E = v_1 p$ is fitted to the point at $p/\hbar = 0.55\,A^{-1}$, a value of 239 ± 5 meters/sec is obtained for v_1. The measured spectrum passes through a maximum of $E/k = 13.92 \pm 0.10$°K at $p/\hbar = 1.11 \pm 0.02\,A^{-1}$. There is a minimum at $p/\hbar = 1.92 \pm 0.01\,A^{-1}$, which may be represented by Landau's roton expression, $E = \Delta + (p - p_0)^2/2\mu$, with $\Delta/k = 8.65 \pm 0.04$°K, $p_0/\hbar = 1.92 \pm 0.01\,A^{-1}$, $\mu = (0.16 \pm 0.01)\, m_{He}$. Above $p/\hbar = 2.18\,A^{-1}$, the spectrum rises linearly with a

slope corresponding to the velocity of first sound. Data were obtained in the region of the minimum at $T = 1.6°K$ and at $T = 1.8°K$. The spectrum has the same general shape observed at $T = 1.1°K$, shifted downward by 0.22°K at 1.6°K, and by 0.50°K at 1.8°K. The temperature variation may be represented by the empirical expression $\Delta/k = 8.68 - 0.0084T^7°K$. At the higher temperatures, the energy spread of the excitations was larger than the energy resolution of the apparatus. The observed full width at half maximum was ~1°K at $T = 1.6°K$, and ~2°K at $T = 1.8°K$.

EXCITATIONS IN LIQUID HELIUM: THERMODYNAMIC CALCULATIONS. P. J. Bendt, R. D. Cowan, and J. L. Yarnell, University of California, Los Alamos Scientific Laboratory, Los Alamos, New Mexico (Received October 16, 1958).

The entropy, specific heat, normal fluid density, and velocity of second sound in liquid helium II have been calculated by applying statistical mechanics to the thermal excitations. The calculations were based on the energy-momentum relation obtained by neutron scattering measurements described by Yarnell, Arnold, Bendt, and Kerr, and were made on an IBM-704 electronic digital computer by numerical integrations over the observed excitation curve. A better approximation than Landau's has been obtained by extending Landau's theory to take account of the temperature dependence of the excitation curve. An expression of the form $E(p,T) = c - d(\rho_n/\rho)$ was used to interpolate the excitation energy between temperatures at which it was measured. Results between 0.2 and 1.8°K are not sensitive to the exact form of the interpolation expression. Agreement of the calculations with experimental measurements is as follows: Entropy, ±3% in the temperature range 0.2 to 1.8°K; specific heat, ±4% between 0.2 and 1.7°K; second sound velocity, ±4% between 0.8 and 1.8°K, and ±2% between 1.0 and 1.7°K. The calculated normal fluid density ρ_n agrees with experimental values derived from second sound velocity and specific heat measurements within ±8% between 0.7 and 2.0°K and within ±5% from 1.1 to 1.9°K. These values are, however, higher than those obtained from torsion pendulum measurements, which are 27% below the calculated value at 1.2°K. Also calculated as functions of temperature are the average effective mass (as defined by Landau) of excitations in four momentum intervals, and values of $-\kappa/\overline{B}$, the thermal conductivity κ divided by the average over momentum of the Khalatnikov nonequilibrium kinetic coefficient $-B$, and $\eta/\overline{C}$, the viscosity η divided by the average value of the Khalatnikov coefficient C.

THERMODYNAMICS OF INHOMOGENEOUS SYSTEMS. II. Edward W. Hart, General Electric Research Laboratory, Schenectady, New York (Received November 7, 1958).

By the inclusion of an explicit dependence of the energy on the boundary of the system, the thermodynamic description of inhomogeneous systems is made more complete. In particular, virtual processes can then be defined that determine the measurable stresses everywhere in the system.

THIRD AND FOURTH SOUND IN LIQUID HELIUM II. K. R. Atkins, Department of Physics, University of Pennsylvania, Philadelphia, Pennsylvania (Received October 29, 1958).

This article discusses the possible existence of two hitherto undetected types of wave propagation in liquid helium II. Third sound is a surface wave of long wavelength on a liquid helium film during which the normal component remains stationary and the superfluid component oscillates parallel to the wall. To treat this properly it is necessary to consider temperature changes and evaporation from the surface of the film. Fourth sound may exist in narrow two-sided channels. The normal component again remains stationary and the superfluid component oscillates parallel to the wall, but the width of the channel must remain fixed and so there are oscillations in both total density and temperature.

CLASSICAL FLUCTUATION-RELAXATION THEOREM. Robert H. Kraichnan, Institute of Mathematical Sciences, New York University, New York, New York (Received September 29, 1958).

A general expression is derived for the average infinitesimal-impulse-response matrix of a conservative classical system in a canonical ensemble. The equations of motion are taken as $\dot{x}_n$

$= X_n$, where the X's, as well as the energy E, are functions of the x's but not of their time derivatives. The Liouville equation $\sum \partial \dot{x}_n / \partial x_n = 0$ is assumed, but it is not required that the equations of motion be derivable from a Hamiltonian. If they are, the x's are the canonical coordinates and momenta. The result found is $g_{mn}(\tau) = \Phi_{mn}(\tau)/kT$, $(\tau \geq 0)$, where $\epsilon g_{mn}(\tau)$ is the average increment in x_m at time t resulting from an infinitesimal increment ϵ externally induced in x_n at time $t-\tau$; $\Phi_{mn}(\tau)$ is the covariance $\langle x_m(t)\,(\partial E/\partial x_n)' \rangle$, where the prime denotes argument $t-\tau$; k and T are Boltzmann's constant and absolute temperature. This relation is derived as a direct consequence of the fact that two initially isolated systems in equilibrium at identical temperatures remain in equilibrium when weakly coupled to each other.

PHOTOCONDUCTOR-PERFORMANCE, SPACE-CHARGE CURRENTS AND THE STEADY-STATE FERMI LEVEL. A. Rose and M. A. Lampert, RCA Laboratories, Princeton, New Jersey (Received September 29, 1958).

The performance of a photoconductor is analyzed, via the concept of the steady-state Fermi level, and shown to be limited by the injection of space-charge. Using the gain-bandwidth product G/τ_0 as a measure of performance, it is found that $G/\tau_0 = M/\tau_\gamma$, where τ_γ is the dielectric relaxation time under operating conditions, and $M = \mathfrak{N}_A/\mathfrak{N}_T$, with $e\mathfrak{N}_A$ the total charge on the anode and $e\mathfrak{N}_T$ the total volume charge, free plus trapped, effectively in thermal contact with the free charge. Generally $M>1$ is achieved only concomitantly with space-charge-limited currents varying steeply with voltage. An important exception is the case where recombination centers control the onset of injected space charge.

TRANSIENT BEHAVIOR OF THE OHMIC CONTACT. M. A. Lampert and A. Rose, RCA Laboratories, Princeton, New Jersey (Received September 29, 1958).

Under certain conditions the response time of a solid for current transients may be determined by the properties of the injecting contact rather than of the bulk. In this case it is the time required for readjustment of the space-charge barrier at the contact for delivery of the new steady-state current. The critical parameter determining the response time of the contact is the total amount of excess charge within one Debye length of the potential minimum in the space-charge barrier. Trapped, as well as free, excess charge must be included in the calculation of the Debye length. The product of gain (G) and bandwidth ($1/\tau_0$) for a photoconductor whose response time is contact-controlled is derived and expressed in the "universal" form $G/\tau_0 = M/\tau_\gamma$, where τ_γ is the dielectric relaxation time under operating conditions.

ULTRASONIC ATTENUATION BY ELECTRONS IN METALS. Eugene I. Blount,* Physics Department and Institute for the Study of Metals, University of Chicago, Chicago, Illinois (Received April 7, 1958).

The calculation of ultrasonic attenuation is discussed for arbitrary frequency and band structure. The results are applied to the cases of a metal with a spherical energy surface and a semimetal with electrons in two "valleys." Possible experimental uses are pointed out. The occurence of saturation is discussed. The acoustoelectric effect is also treated by similar methods and is seen to be closely related to the attenuation.

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INFRARED ABSORPTION OF REDUCED RUTILE TiO_2 SINGLE CRYSTALS. D. C. Cronemeyer, Bendix Aviation Corporation, Research Laboratories, Detroit, Michigan (Received May 7, 1958).

Rutile TiO_2 single crystal plates have been reduced in hydrogen at about 700°C for several minutes to make them semiconducting. The concentration of oxygen vacancies was controlled by variations of time and temperature. The infrared absorption of a series of plane parallel plates having electrical resistivities ranging from 3 to 0.01 ohm-m has been examined. It is postulated that the electrical conductivity arises from the ionization of either one or two trapped electrons from each oxygen vacancy.

In samples with electrical resistivity ($\perp$ to the c axis) greater than 0.04 ohm-m, the optical absorption at room temperature peaks at about 0.75 ev. For samples with electrical resistivity

less than 0.03 ohm-m, the optical absorption shows a new maximum at 1.18 ev. The decrease of thermal activation energy with increasing oxygen vacancy concentration is expected to explain the "optical transition" from 0.75 to 1.18 ev. The ionization energies agree reasonably well with those calculated for a helium atom model of a doubly ionizable donor immersed in a dielectric medium [$K_\varepsilon = n_0^2 = (2.40)^2$], namely 0.73 ev and 1.64 ev. A modification of this theory is also indicated which predicts the second ionization energy as 1.41 ev in better agreement with the experimental value of 1.18 ev.

STUDY OF FERROELECTRIC TRANSITIONS OF SOLID-SOLUTION SINGLE CRYSTALS OF $KNbO_3$-$KTaO_3$. S. Triebwasser, International Business Machines Corporation, Poughkeepsie, New York (Received November 13, 1958).

Methods of preparation and results of measurements of the dielectric constant of solid-solution single crystals of $KNbO_3$ and $KTaO_3$ are presented. Data are given for the composition range 0-80 mole percent $KTaO_3$ in the temperature interval -180°C to 450°C. It is found that all of the crystals studied show Curie-law behavior of the dielectric constant in the paraelectric state. The data indicate that the paraelectric to ferroelectric transition which is first order for pure $KNbO_3$ becomes second order at a concentration of approximately 55 mole percent $KTaO_3$. Analysis of the results in terms of Slater's model of ferroelectricity in a perovskite lattice leads to the conclusion that the dependence of the Lorentz correction on volume plays a major role in determining the ferroelectric behavior. Ionic polarizabilities and their dependence on volume are also discussed.

ACOUSTOELECTRIC EFFECT IN n-TYPE GERMANIUM. Gabriel Weinreich, T. M. Sanders, Jr.,* and Harry G. White, Bell Telephone Laboratories, Murray Hill, New Jersey (Received November 13, 1958).

Theory and experiment are presented for the drag exerted on electrons in a solid by a traveling ultrasonic wave. After a discussion of the reasons why the effect is generally very small, it is shown that under certain conditions in n-type germanium it may be quite appreciable, and that its size directly indicates the intervalley scattering rate. Experimental data are given for arsenic-doped germanium ranging in impurity content from 10^{14} to 10^{16} cm^{-3}, at temperatures from 20 to 160°K. These yield the absolute value of the "uniaxial" deformation potential constant (16 ev) and the intervalley scattering rate as a function of temperature and doping. The interpretation of the results ascribes intervalley scattering action both to phonons and to impurities. The phonon contribution yields the frequency of the ⟨100⟩ longitudinal phonon (6.6×10^{12} sec^{-1}) and the size of the appropriate coupling parameter; the impurity contribution yields intervalley scattering cross sections for neutral and ionized donors as a function of temperature. The two cross sections are explained as being due, respectively, to exchange scattering events and to compound capture-reemission processes. Further analysis gives a lower limit to the valley-orbit splitting of the arsenic donor ground state (1.7×10^{-3} ev). The experiment verifies that the conduction-band valleys lie on ⟨111⟩ axes at the Brillouin zone edge. Ultrasonic attenuation due to intervalley scattering is discussed and shown to be too small to be easily measurable.

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PHOTOELECTRIC EFFECTS IN CERTAIN OF THE ALKALI HALIDES IN THE VACUUM ULTRAVIOLET. John W. Taylor* and Paul L. Hartman, Physics Department, Cornell University, Ithaca, New York (Received October 7, 1958).

An extensive search for internal photoconductivity in NaCl, KCl, KI, and LiF resulting from stimulation by light of wavelength lying below the fundamental absorption edges of these crystals has yielded negative results. Photon energies employed ranged from about 6 ev to 21 ev. Calculations based on observed external photoelectric emission currents indicate that photoelectric quantum efficiencies as low as 10^{-8} electron per incident photon could have been detected if electrons traveled completely through the crystal.

Measurements of the quantum efficiencies for external photoelectric emission by some of these crystals support the conclusion of Phillip and Taft that the valence bands of CsI and KI lie about 6.0

and 7.2 ev below the vacuum potential, respectively, and indicate that the valence bands of NaCl and KCl are 8.3±0.1 ev below the vacuum. The valence band in LiF is tentatively placed at about 13 ev below the vacuum level.

Energy distribution curves for the photoelectrons ejected from NaCl and KCl by photons of energies of from 10 to 21 ev show behavior similar to that reported for Cs_3Sb, indicating that photoelectrons which have sufficient kinetic energy in the crystal to ionize another electron from the valence band will lose relatively little energy in escaping.

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INFRARED ABSORPTION AND ELECTRON EFFECTIVE MASS IN *n*-TYPE GALLIUM ARSENIDE. W. G. Spitzer and J. M. Whelan, Bell Telephone Laboratories, Incorporated, Murry Hill, New Jersey (Received November 4, 1958).

The infrared absorption between 0.85 and 25 microns has been measured as a function of carrier concentration for *n*-type single crystal gallium arsenide. The absorption in the 1- to 5-micron region is compatible with a model in which there are minima ~0.25 ev above the bottom of the conduction band. Infrared reflectivity measurements on several samples of different carrier concentrations were used to deduce the free-carrier contribution to the electric susceptibility and the electron effective mass. The results indicate a value for the mass of $(0.078 \pm 0.004)m$ with an indication of an increase for the sample of highest carrier concentration. This value is substantially larger than previously reported values.

ACOUSTICAL ABSORPTION ARISING FROM MOLECULAR RESONANCE IN SOLIDS. Leonard Liebermann, University of California, La Jolla, California (Received May 26, 1958).

Resonance phenomena resulting in anomalously high acoustic absorption can occur whenever lattice and internal molecular vibrational frequencies overlap. Significant effects are predicted at all acoustic frequencies, without restriction to the resonance region. A method is given for calculating resonance absorption using available molecular constants. As an illustrative example, the magnitude of resonance absorption has been calculated for benzene. The predicted absorption is 0.5 cm^{-1} at the 10-Mc/sec frequency, orders of magnitude larger than has been observed in other single crystals. In order to test this prediction, measurements of absorption were made on a large single crystal of benzene yielding at 10 Mc/sec, 0.25 cm^{-1}. This agreement between absorption and calculated values is considered adequate, confirming the phenomenon.

SUPERCONDUCTIVITY OF In_2Bi. R. E. Jones and W. B. Ittner, Research Laboratory, International Business Machines Corporation, Poughkeepsie, New York (Received November 4, 1958).

In_2Bi is found to be superconducting at 5.6±0.1°K. The compound InBi is not superconducting above 1.5°K.

ELECTRON INTERACTION IN SOLIDS, CHARACTERISTIC ENERGY LOSS SPECTRUM. P. Nozières, Laboratoire de Physique, École Normale Supérieure, Paris, France, and D. Pines, Institute for Advanced Study, Princeton, New Jersey (Received October 23, 1958).

The characteristic energy loss spectrum of solids is analyzed with the aid of the dielectric formulation of the many-body problem developed by the authors. It is shown that a measurement of the angular distribution of inelastically scattered fast electrons is a direct measurement of the imaginary part of the inverse dielectric constant of the solid, at the frequency and momentum of the energy transfer to the electrons in the solid. The calculations of the dielectric constant by means of first-order perturbation theory and within the random phase approximation are discussed with regard to their implications for the characteristic energy loss spectrum. It is shown that plasmons provide the dominant excitation mode at low momentum transfer for the free electron gas and for most solids; an explicit calculation is given of the role of screening in reducing the role of the independent electron excitations in this region. A comparison is made between characteristic energy loss experiments and other techniques for exploring the energy level structure of solids, including the nonrelativistic Compton effect, optical transmission and reflection experiments, and x-ray emission and absorption studies.

An analogous formulation of the optical properties of solids is given, which shows that in general the transverse and longitudinal dielectric constants will not be equal. They are equal within the random phase approximation, so that a careful comparison of optical data with characteristic energy loss data permits a direct verification of the validity of the random phase approximation. Such a comparison likewise yields information on the magnitude of the screening of individual electron transitions and permits the unambiguous identification of plasmon excitations. Finally it is shown that to the extent that the random phase approximation is valid, the results of a characteristic energy loss experiment may be used for a direct experimental determination of the ground-state energy of the electrons in a solid.

TRANSITIONS FROM FERROMAGNETISM TO ANTIFERROMAGNETISM IN IRON-ALUMINUM ALLOYS: EXPERIMENTAL RESULTS. Anthony Arrott and Hiroshi Sato, Scientific Laboratory, Ford Motor Company, Dearborn, Michigan (Received October 16, 1958).

The magnetic behavior of iron-aluminum alloys has been determined for compositions from 0 to 40 atomic percent aluminum in external fields to 12 kilo-oersteds from 4°K to 300°K. The alloys above 33 atomic percent exhibit antiferromagnetism. Those below 28 atomic percent are ferromagnetic. Between 28 and 33 atomic percent, the behavior is ferromagnetic at room temperature but shows a transition to antiferromagnetism on lowering the temperature. The effect of order on the magnetic properties has also been investigated. Antiferromagnetism is found only in the ordered structures.

TRANSITIONS FROM FERROMAGNETISM TO ANTIFERROMAGNETISM IN IRON-ALUMINUM ALLOYS: THEORETICAL INTERPRETATION. Hiroshi Sato and Anthony Arrott, Scientific Laboratory, Ford Motor Company, Dearborn, Michigan (Received October 16, 1958).

Complicated magnetic properties of α-phase iron-aluminum alloys are successfully described from a theory with a few simple assumptions concerning the magnetic interaction between iron atoms. Two interactions are considered: the direct exchange interaction between nearest neighboring iron atoms and the indirect super-exchange interaction between iron atoms separated by an aluminum atom. The model also leads to a reinterpretation of the neutron diffraction data. In view of the success of this model, it is to be concluded that the direct exchange interaction between nearest neighbor iron atoms is mainly responsible for the ferromagnetism of α-iron.

POTENTIAL FLUCTUATIONS IN THE TRANSITION REGION TO SUPERCONDUCTIVITY. Hans Meissner, Johns Hopkins University, Baltimore, Maryland (Received September 30, 1958).

Gorter recently proposed a dynamical model to describe the state of current-carrying wire, in the transition region to superconductivity. This model predicts in the transition region to superconductivity an alternating component of the potential difference across a wire carrying a strong direct current. The magnitude of this fluctuation should be of the order of a few percent of the average potential. Using a shielded transformer in the liquid helium and a battery-operated preamplifier, detection of this alternating component has been attempted. The noise output was somewhat larger when the sample was in the intermediate state than when it was in the superconducting state. However, instead of being a few percent of the average potential (as in the Gorter model), this additional alternating potential was only about one hundredth of one percent. It may be due partially to a slightly drifting bath temperature, since vibrations made it necessary to switch the pump off about 10 seconds before the readings were taken.

NUCLEAR RELAXATION IN ALUMINUM. John J. Spokas* and Charles P. Slichter, Department of Physics, University of Illinois, Urbana, Illinois (Received November 3, 1958).

Nuclear spin lattice relaxation times T_1 and phase memory times T_2 have been measured in 99.99% pure aluminum foil from 77°K to the melting point by the spin-echo method. At high temperature the line narrows due to self-diffusion, as previously observed by Seymour. The temperature dependence of the narrowing gives an activation energy of self-diffusion E_D of 1.4 ± 0.1 ev. From the measured "narrowing temperature"

and Nowick's theoretical estimate of D_0, the same value of E_D is found. As predicted theoretically, measurements show that T_1 arises from interaction with conduction electrons and is inversely proportional to temperature from 1°K to 930°K. At the highest temperatures, an unknown line-broadening mechanism is found which is believed to represent quadrupole coupling to strain fields of dislocations. The measurements were made feasible only through the use of a phase-coherent detection system, a description of which is given.

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ANTIFERROMAGNETIC MAGNON DISPERSION LAW AND BLOCH WALL ENERGIES IN FERROMAGNETS AND ANTIFERROMAGNETS. R. Orbach, Department of Physics, University of California, Berkeley, California (Received October 16, 1958).

The exact eigenstates of the exchange Hamiltonian, $\mathcal{H} = 2J\sum_l(\vec{S}_l \cdot \vec{S}_{l+1} - \frac{1}{4})$, are found for short chains of 4, 6, 8, and 10 atoms of spin 1/2. A linear dispersion law for magnons in an antiferromagnet is exhibited by the energy spectrum. The periodic boundary conditions are then removed and the ends of the chain held fixed, both parallel and antiparallel. The energy of the 180° Bloch wall is computed and compared with the classical result. It is found that the semiclassical ferromagnetic wall is a good approximation to the exact wall. The energy of the semiclassical antiferromagnetic wall is not a very good approximation to the exact wall energy, but the semiclassical energy appears to have the correct dependence on the wall thickness.

EXCLUDED VOLUME PROBLEM AND THE ISING MODEL OF FERROMAGNETISM. Michael E. Fisher and Martin F. Sykes, Wheatstone Physics Laboratory, King's College, London, England (Received November 3, 1958).

The relationship between the excluded volume problem for a discrete random walk on a lattice and the corresponding Ising model of ferromagnetism is investigated. Systematic methods are presented for the construction of rigorous lower bounds to the limit $\mu = \lim_{n\to\infty}(c_{n+1}/c_n)$, where c_n is the number of n-step self-avoiding walks on a given lattice. In this way Temperley's conjecture that $\mu = \coth(J/kT_C)$, where T_C is the Curie temperature of the corresponding Ising-model ferromagnet, is disproved. The series c_n for various two- and three-dimensional lattices have been enumerated exactly for values of n from ten to twenty. Extrapolation of these series by procedures known to be valid from exact Ising-model results, yields more accurate values of μ than Wall's statistical calculations and also shows that $c_n \sim n^{\alpha}\mu^n$, where $\alpha \approx 1/3$ for plane lattices and $\alpha \approx 1/7$ for three-dimensional lattices. This means that the entropy of the nth "link" of a polymer molecule in solution should vary as $\delta S_n = k\ln\mu + k\alpha/n$. The relevance of these results to the interpretation of the boundary tension of the Ising model, to the critical behavior of gases, and to the mean square size of a polymer molecule is discussed briefly.

RADIATIVE ALPHA PARTICLE CAPTURE IN Ni^{58} AND Xe^{136} AT HIGHER ENERGIES. J. B. Ball, A. W. Fairhall, and I. Halpern, Departments of Chemistry and Physics, University of Washington, Seattle, Washington (Received November 10, 1958).

The cross sections of the (α,γ) reaction have been measured as a function of bombarding energy in Ni^{58} and Xe^{136} up to about 35 Mev. In both cases the capture events were identified by means of the specific radioactivity produced in the reaction. In the bombardment of Ni^{58} it was necessary to make use of the tight angular spread of recoils in the capture reaction to distinguish these recoils from those of the $Ni^{60}(\alpha, 2n)$ reaction. Both measured capture cross sections rise to a maximum value (Ni^{58}: 400 μb, Xe^{136}: 60 μb) at an excitation energy of about 20 Mev. A comparison of the observed cross sections with those obtained in (p,γ) reactions at similar energies suggests that these higher energy capture reactions take place largely through the formation of a compound nucleus rather than by any "direct" mechanism.

SEARCH FOR RESONANCE STRUCTURE OF NEUTRON CROSS SECTIONS AT 100 Mev. Stuart G. Carpenter* and Richard Wilson, Harvard University, Cambridge, Massachusetts (Received November 10, 1958).

We have searched for resonances in neutron cross sections near 100 Mev by a self-absorption

method. We find a sign, barely significant, of resonance structure in antimony or cadmium, but if the resonance cross section is 1.2 times the nonresonant cross section, the resonances are present over only 3% of the energy band. The upper limit for the other elements is 1%.

We have measured precisely differences of neutron total cross sections for several adjacent elements and find that there are many deviations from a smooth curve as a function of $A^{2/3}$.

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CLOUD CHAMBER MEASUREMENT OF THE HALF-LIFE OF THE NEUTRON. N. D'Angelo, Argonne National Laboratory, Lemont, Illinois (Received November 3, 1958).

The half-life of the free neutron has been measured by a cloud chamber method. It turned out to be 12.7 min with an error of ~ 15%, ~ 10% being due to the poor statistics. This determination has about the same accuracy as the best previous measurement and confirms it by an independent method.

STUDY OF THE NEUTRON REACTIONS $Li^6(n, \alpha)H^3$, $F^{19}(n, \gamma)F^{20}$, and $I^{127}(n, \gamma)I^{128}$. Fletcher Gabbard, Robert H. Davis,* and T. W. Bonner, The Rice Institute, Houston, Texas (Received November 3, 1958).

The reactions $Li^6(n, \alpha)H^3$, $F^{19}(n, \gamma)F^{20}$, and $I^{127}(n, \gamma)I^{128}$ were studied using scintillating crystals both as targets and as detectors. The neutrons were produced by the $Li^7(p, n)Be^7$ and $T^3(p, n)He^3$ reactions. A redetermination of the $Li^7(p, n)Be^7$ differential cross section was made in the proton energy range 1.9 to 3.0 Mev at 0° and the cross section observed at 2.27 Mev was 154 mb. Cross sections for this reaction were determined at 120° for proton energies below 2.4 Mev and at 60° for energies below 2.1 Mev. The $Li^6(n, \alpha)H^3$ reaction cross section was studied for neutron energies between 0.025 and 4 Mev. The resonance at 0.255 Mev was observed to have a width (lab) of 0.125 Mev and a maximum cross section of 2.80 barns. A fit of the dispersion theory to this resonance yielded the following parameters in the center-of-mass coordinate system: $\gamma_n^2 = 0.41$ Mev-cm; $\gamma_\alpha^2 = 0.015$ Mev-cm; $\Delta_\lambda = 0.17$ Mev; $E_\lambda = 0.039$ Mev; and $\Gamma = 0.188$ Mev at the resonant energy. Sixteen resonances were observed in the $F^{19}(n, \gamma)F^{20}$ reaction at neutron energies of 0.027, 0.048, 0.100, 0.177, 0.270, 0.308, 0.388, 0.425, 0.500, 0.600, 0.760, 0.865, 0.950, 1.125, 1.290, and 1.635 Mev. The radiative widths Γ_γ of twelve of these resonances vary between 1.1 and 8.6 ev. A monotonically decreasing cross section was observed for the $I^{127}(n, \gamma)I^{128}$ reaction between 0.025 and 0.500 Mev.

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TOTAL CROSS SECTION FOR n-p SCATTERING AT 20 Mev. R. B. Day, R. L. Mills, J. E. Perry, Jr., and F. Scherb,* University of California, Los Alamos Scientific Laboratory, Los Alamos, New Mexico (Received November 3, 1958).

The total scattering cross section for neutrons incident on protons has been measured in a good-geometry transmission experiment. The scatterer was hydrogen gas contained in a pressure vessel one meter long, at a pressure of 4000 psi. The resulting cross section is $\sigma_t = 0.4942 \pm 0.0025$ barn at a laboratory neutron energy of 19.665 ±0.026 Mev.

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HALOGEN DELAYED-NEUTRON ACTIVITIES. Gilbert J. Perlow and Andrew F. Stehney, Argonne National Laboratory, Lemont, Illinois (Received August 8, 1958).

A gas-flow technique has been developed to make rapid chemical separations of bromine and of iodine fission products from a solution of U^{235} irradiated with thermal neutrons. The active bromine or iodine is observed by its delayed neutron emission. Analysis of the decay curves so obtained has been made by graphical and by computer methods. Four delayed-neutron periods have been found in the bromine fraction and three in the iodine. Including several already well known, the half-lives and probable mass assignments are: 54.5-sec Br^{87}, 16.3-sec Br^{88}, 4.4-sec Br^{89}, and 1.6-sec Br^{90}; 24.4-sec I^{137}, 6.3-sec I^{138}, and 2.0-sec I^{139}. Shorter-lived halogens would not have been detected. The re-

lative yields of each neutron activity in the order of decreasing half-life are, for bromine, 0.37: 1.0: 1.9: 1.5; while for iodine they are 1.0: 0.47: 0.38 in the same order. If we assume that there are no other important contributors of delayed neutrons in this half-life range, the numerical values of yields for bromine may be compared directly with those for iodine with an uncertainty of a factor of 2.

TOTAL NEUTRON YIELDS FROM LIGHT ELEMENTS UNDER PROTON AND ALPHA BOMBARDMENT. J. H. Gibbons and R. L. Macklin, Oak Ridge National Laboratory, Oak Ridge, Tennessee (Received November 10, 1958).

Total neutron production cross sections have been measured for targets of D, T, Li^7, Be^9, B^{11}, C^{13}, C^{14}, and F^{19} under proton bombardment and of Li^7, Be^9, B^{10}, and Si^{29} under alpha particle bombardment. Energy ranges covered differ but were contained between threshold and 5 Mev for protons and between threshold and 9 Mev for alphas. In several cases cross sections for the inverse reactions (n, p) and (n, α) have been computed and compared, where available, with direct measurements.

EFFECT OF CONFIGURATION MIXING ON MAGNETIC MULTIPOLE RADIATION. K. K. Gupta* and R. D. Lawson, Enrico Fermi Institute for Nuclear Studies, University of Chicago, Chicago, Illinois (Received November 12, 1958).

Expressions for the matrix element governing the emission of magnetic multipole radiation are given. The effect of weak configuration mixing and the small additional contribution to proton matrix elements (~5%) arising from the nuclear spin-orbit interaction are considered. For $M4$ transitions, even in nuclei near closed shells, the theoretical matrix elements derived from pure configurations are too large by a factor of two. The effect of configuration mixing reduces these. Choosing a δ-function interaction of reasonable strength leads to a reduction of ~25% in the matrix element for the $1g_{9/2} \rightarrow 2p_{1/2}$ transition in $_{39}Y^{89}$. The $(2p_{1/2})^{-1} \rightarrow (1g_{9/2})^{-1}$ transition in $_{38}Sr^{87}$ is reduced close to the required 50%. The $1h_{11/2} \rightarrow 2d_{3/2}$ matrix element becomes ~25% smaller. However, there are no data available for the nucleus $_{50}Sn^{115}$, to which one would like to compare. Finally, the $(1i_{13/2})^{-1} \rightarrow (2f_{5/2})^{-1}$ transition in $_{82}Pb^{207}$ has its matrix element reduced by only ~6%. Although the reduction in $_{39}Y^{89}$ is somewhat too small, one cannot make a positive statement that configuration mixing is incapable of explaining the experimental result since a two-body force of a more realistic type may close the gap between theory and experiment. However, for $_{82}Pb^{207}$ there seems no doubt that other effects, for example, meson currents, must be incorporated.

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LIGHT-PARTICLE SPECTRA FROM THE NITROGEN BOMBARDMENT OF OXYGEN. M. L. Halbert and A. Zucker, Oak Ridge National Laboratory, Oak Ridge, Tennessee (Received November 10, 1958).

Energy spectra and angular distributions were measured for protons and α-particles resulting from the nitrogen bombardment of oxygen. It was found that the angular distributions of protons are reasonably isotropic whereas α-particles are peaked in the forward direction. The energy distributions at various angles were compared with the predictions of a simple level density formula $w = C \exp[2(aE^*)^{1/2}]$. The constant a was found to be $a = 2.5 \pm 0.2$ from proton spectra, $a = 1.3 \pm 0.2$ from α-particle spectra at all angles measured. The energy spectrum of deuterons was measured at 15 deg and a value of $a = 1.8 \pm 0.3$ was found. Unexplained structure was observed in the spectrum of α-particles emitted at zero degrees. After correction for extraneous factors, the results indicate that α-particles are emitted on the order of ten times more copiously than protons, while deuterons are about half as abundant as protons.

ZEEMAN INVESTIGATIONS OF CURIUM-242. J. C. Hubbs, R. Marrus, and J. O. Winocur, Radiation Laboratory and Department of Physics, University of California, Berkeley, California (Received November 3, 1958).

The atomic-beam magnetic-resonance method has been used to investigate 163-day Cm^{242}. The spin of this even-even nuclide is found to be zero. Four low-lying electronic energy levels

are found and the Lande g factors are measured to be $g_{J_2}=2.561\pm0.003$, $g_{J_3}=2.000\pm0.003$, $g_{J_4}=1.776\pm0.002$, and $g_{J_5}=1.671\pm0.003$. No direct measurement can be made of the angular momenta of these levels, but other considerations contained in the text make highly probable the J values indicated in subscripts and arising from the electronic configuration $(5f)^7(6d)^1(7s)^2$.

DISTRIBUTION OF THE ANGULAR MOMENTA, LEVEL SPACINGS, AND NEUTRON WIDTHS OF Al^{28}. Carl T. Hibdon, Argonne National Laboratory, Lemont, Illinois (Received November 7, 1958).

The neutron total cross section from about 1 kev to 450 kev shows the presence of sixty-six peaks, seven of which are of the s-wave variety. Thirteen resonances are attributable to $J=0$, twenty-one to $J=1$, eighteen to $J=2$, ten to $J=3$ and three to $J=4$. This distribution is in agreement with the theoretical distribution for a value of $2c\tau\equiv 2\sigma^2=6$. The density of all levels for this energy interval is 146 Mev^{-1} and the average level spacing of the nucleons is 0.48 Mev. The neutron widths vary from 1 to 7 kev and the distribution of the reduced widths appears to agree with an exponential distribution and is also in fair agreement with the Porter-Thomas distribution. The level spacings also agree with an exponential distribution. As obtained from the reduced widths averaged over both values of J, the value of the strength function for $l=0$ is 0.05, averaged over all values of J for $l=1$ it is 0.49, and for $l=2$ it is too large in comparison with the p-wave strength function. The particularly low value of the cross section below 30 kev and the shape of the wings of the 35-kev resonance can be explained by a multiple-level computation of the interference of the s-wave levels. On the basis of the results of the present analysis, the levels are about equally divided between the odd and even values of l.

T(p,p)T SCATTERING NEAR THE T$(p,n)$$He^3$ THRESHOLD. Nelson Jarmie and Robert C. Allen,* Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico (Received November 7, 1958).

The absolute cross section of T(p,p)T elastic scattering at center-of-mass angles of 58° 38′ and 109° 31′ has been measured as a function of energy near the T$(p,n)$$He^3$ threshold. Small but distinct perturbations are observed near the threshold and are discussed in the light of theoretical predictions. The bearing of the data on a possible excited state in He^4 at 20 Mev is discussed.

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DECAY PROPERTIES OF I^{136}. Noah R. Johnson and G. D. O'Kelley, Oak Ridge National Laboratory, Oak Ridge, Tennessee (Received November 10, 1958).

I^{136}, produced from the slow neutron fission of uranium, has been studied with scintillation techniques. From the single-crystal (NaI) and gamma-gamma coincidence spectra, gamma rays were found at energies of 0.20, 0.27, 0.39, 0.46, 0.53, 0.71, 1.00, 1.321±0.006, 1.55, 1.72, 1.91, 2.25, 2.40, 2.63, 2.84, and 3.2 Mev. Also, from single-crystal (anthracene) and beta-gamma coincidence measurements, four beta-ray groups were observed with energies of 2.73, 4.23, 5.62, and 7.00 Mev, where the latter represents the I^{136}-Xe^{136} mass difference and is present in about 6% of the total beta-ray transitions. A partial decay scheme has been proposed. The half-life of I^{136} was measured as 82.8±1.5 seconds.

NUCLEAR ENERGY LEVELS AND MASS OF NITROGEN-17. C. S. Littlejohn, Enrico Fermi Institute and Department of Physics, University of Chicago, Chicago, Illinois (Received November 7, 1958).

Energy levels and the atomic mass defect of the nucleus N^{17} have been determined by measuring the kinetic energies of proton groups produced in the reaction $B^{11}(Li^7,p)N^{17}$ on bombardment of isotopically separated B^{11} targets with 2-Mev Li^7 ions.

The protons were counted at 90° in the laboratory system from the forward direction of the lithium beam. The group of highest energy gave an $(M-A)$ value of 12.92±0.06 Mev for the N^{17} ground state, corresponding to a physical atomic weight of 17.01388±0.00006. Two completely resolved groups of lower energy showed excited states at 1.32 and 1.89 Mev. Two strong but in-

completely separated peaks showed levels at 2.50 and 2.82 Mev of excitation. Incompletely resolved groups of lower energy indicated four higher states at 3.27, 3.57, 3.86, and 4.18 Mev. Due to low counting rates, intensity estimates are difficult but indicate that, at 90° and 2 Mev bombardment, the ground state group is only 1/4 as strong as the most intense groups which arise from the 2.50- and 2.82-Mev levels. The groups from the higher levels (3.3 - 4.2 Mev) are less intense by a factor of 0.7. A comparison is made of the spectrum of N^{17} with O^{17} and F^{17}; Coulomb corrections are applied to calculate the isotopic spin $T=3/2$ sequence corresponding to the above levels in O^{17}, the lowest such level being predicted at 11.04 Mev in O^{17}. A brief discussion is given of the reaction and also of the interpretation in terms of nuclear shell structure of the levels observed.

CROSS SECTIONS FOR NUCLEAR REACTIONS INVOLVING NUCLEAR ISOMERS. Bruno Linder* and Ralph A. James, Department of Chemistry, University of California, Los Angeles, California (Received November 3, 1958).

Cross sections for the reactions $Cr^{52}(p,n)Mn^{52m}$, $Cr^{53}(p,n)Mn^{53}$, $Ag^{107}(p,pn)Ag^{106m}$, $Ag^{107}(p,pn)Ag^{106}$, and $Cu^{63}(p,n)Zn^{63}$ were determined with the 20.6-Mev proton beam of the U.C.L.A. synchrocyclotron as a function of energy. The relative yields of Sr^{85}, Sr^{85m}, and Sr^{87m} obtained in the proton bombardment of rubidium were also determined. In all cases the isomer yield ratios varied with energy over the entire energy range studied. The results are discussed in terms of conservation of angular momentum, assuming compound nucleus formation.

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DISINTEGRATION OF Ag^{110m}. H. W. Taylor and S. A. Scott, Physics Department, Queen's University, Kingston, Canada (Received August 14, 1958; revised manuscript received December 18, 1958).

The gamma radiations of Ag^{110m} have been studied using a conventional scintillation coincidence spectrometer and the gamma-ray summing technique. The positions of the 1484, 735, and 1560 kev gamma-rays in the decay scheme have been determined.

GAMMA-GAMMA DIRECTIONAL CORRELATION IN Cd^{110}. H. W. Taylor and W. R. Frisken, Physics Department, Queen's University, Kingston, Canada (Received August 14, 1958; revised manuscript received December 18, 1958).

Directional correlation experiments have been performed for the 1385-885, 760-1510, and 940-885 kev gamma-ray cascades in Cd^{110}. Spins of 5, 6, 3, 4 have been assigned to the levels at 2935, 2490, 2170, and 1540 kev in the Cd^{110} level scheme. The 1385- and 1510-kev gamma-rays are dipole-quadrupole mixtures.

EXCITATION FUNCTIONS OF BISMUTH AND LEAD. W. J. Ramler, J. Wing, D. J. Henderson, and J. R. Huizenga, Argonne National Laboratory, Lemont, Illinois (Received November 14, 1958).

Cross sections for the reactions $Bi^{209}(He^4,2n)At^{211}$, $Bi^{209}(He^4,3n)At^{210}$, and $Bi^{209}(He^4,4n)At^{209}$ were measured with helium ions of energies from 20.6 to 43.3 Mev. Cross sections for the reactions $Bi^{209}(d,p)RaE$, $Bi^{209}(d,n)Po^{210}$, $Bi^{209}(d,2n)Po^{209}$, $Bi^{209}(d,3n)Po^{208}$, and $Pb^{208}(d,p)Pb^{209}$ were measured with deuteron ions of energies from 6.3 to 21.5 Mev. Some information is presented on the $Bi^{209}(He^4,t)Po^{210}$ reaction. A half-life of 7.23±0.04 hr was measured for At^{211}. The $Bi^{209}(d,3n)Po^{208}$ threshold energy is 12.0±0.3 Mev (center-of-mass system). The compound nucleus reaction cross sections are compared with the predictions of the Jackson model, and good agreement is obtained. The (d,n) and (d,p) stripping cross sections are discussed briefly.

$C^{13}(p,\gamma)N^{14}$ 1.47- AND 2.11-Mev RESONANCES AND THE ODD-PARITY LEVELS OF MASS 14. B. K. Warburton,* Brookhaven National Laboratory, Upton, New York, and Palmer Physical Laboratory, Princeton University, Princeton, New Jersey; and H. J. Rose† and E. N. Hatch,‡ Brookhaven National Laboratory, Upton, New York (Received November 11, 1958).

A study has been made of the γ rays emitted by the 8.9- and 9.5-Mev levels of N^{14} using the $C^{13}(p,\gamma)N^{14}$ reaction at the 1.47- and 2.11-Mev

resonances. The decay schemes of these levels were reinvestigated using a three-crystal pair spectrometer, NaI(Tl) single-crystal measurements, and standard γ-γ coincidence techniques. The anisotropies relative to the bombarding beam of most of the observed γ-transitions were measured using the three-crystal pair spectrometer. The angular distributions of some of the γ-transitions were measured using a single NaI(Tl) crystal. Measurements were made of the Doppler shifts of the ground state decay of the N^{14} 5.10-Mev level and the cascade from the N^{14} 5.83-Mev level to the N^{14} 5.10-Mev level. From these Doppler shift measurements the mean lifetime of the N^{14} 5.83-Mev level was found to be in the range $5\text{-}6.5\times10^{-14}$ sec, while an upper limit of 3×10^{-13} sec was set on the mean lifetime of the N^{14} 5.10-Mev level. The results of the study of the γ-decay of the N^{15} 9.5- and 8.9-Mev levels combined with the results of earlier measurements give conclusive assignments of 2^-, 3^-, 3, and 2 for the N^{14} levels at 9.50, 8.90, 5.83, and 5.10 Mev. The 5.83-Mev level most probably has odd parity. A tentative assignment of $J=2$ is given to the N^{14} 7.02-Mev level. Evidence is presented, from this and previous investigations, that indicates the N^{14} 8.06-, 8.70-, 8.90-, and 9.50-Mev levels arise from the $s^4p^9 2s$ and s^4p^9d configurations with the largest contribution being from the $(p_{1/2}\ 2s_{1/2})$ configuration for the 8.06- and 8.70-Mev levels and from the $(p_{1/2}\ d_{5/2})$ configuration for the 8.90- and 9.50-Mev levels. The analogs of these $T=1$ levels in C^{14} are almost certainly the C^{14} 6.09-, 6.89-, 6.72-, and 7.35-Mev levels, respectively. The $T=0$ $s^4p^9 2s$ and s^4p^9d states of N^{14} are also discussed. It is proposed that the 4.91-, 5.69-, 5.10-, and 5.83-Mev levels of N^{14} are 0^- and 1^- $s^4p^9 2s_{1/2}$ and 2^- and 3^- $s^4p^9d_{5/2}$ states, respectively.

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SCATTERING OF POLARIZED ELECTRONS BY POLARIZED NUCLEONS. James H. Scofield, Indiana University, Bloomington, Indiana (Received November 5, 1958).

The differential cross section for the scattering of arbitrarily polarized charged spin one-half particles by arbitrarily polarized charged spin one-half particles with an anomalous magnetic moment has been calculated, including recoil effects.

NUCLEAR PARAMETERS IN THE SCATTERING OF NUCLEONS BY CARBON. Richard Wilson, Cyclotron Laboratory, Harvard University, Cambridge, Massachusetts (Received November 10, 1958).

Scattering of protons and neutrons by carbon at small angles is discussed. It is shown that it is not possible to obtain a coherent understanding of the scattering at 310 Mev without including an increase of the radius due to nuclear forces. The increase required is in rough agreement with that calculated from nucleon-nucleon scattering.

At 135 Mev, where more complete data are available, it is shown that this increase is not sufficient to obtain a coherent understanding. Either approximations break down or there is a difference in the scattering of neutrons and protons in addition to the Coulomb forces. Further small-angle polarization measurements are badly needed at both energies.

ANGULAR DISTRIBUTIONS AND CORRELATIONS FOR ALPHA-PARTICLE SCATTERING BY C^{12}, Mg^{24}, AND Ca^{40}. G. B. Shook,* University of Washington, Seattle, Washington (Received November 14, 1958).

Angular distributions of alphas scattered from Mg^{24} and Ca^{40} are reported and effective interaction radii are obtained, based on the Born approximation for elastic scattering and on direct-interaction theory for inelastic scattering. For Mg^{24}, levels at 0, 1.37, 4.2, and 6.3 Mev yield interaction radii of 6.4, 6.55, 6.5, and 6.5, respectively (in units of 10^{-13} cm). Evidence is presented in support of the assignment of the 6.3-Mev level to Mg^{24} rather than to Mg^{25} or Mg^{26}, with spin and parity 1-. For Ca^{40}, levels at 0 and 4.48 Mev yield interaction radii of 6.65 and 7.3, respectively. Alpha-gamma angular correlations, which correspond closely to those predicted on basis of direct interaction theory, were obtained for C^{12} (4.43-Mev level) and for Mg^{24} (1.37-Mev level). For Ca^{40} (4.48-Mev level) the correlation corresponds qualitatively to that pre-

dicted for a 1- level with cascade decay through an intermediate state to ground level. Further evidence is presented which supports this interpretation. A higher level in Ca^{40} is observed at 8.4 Mev. All alpha bombardments were made at 43-Mev incident energy.

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ATOMIC MASSES OF THE STABLE ISOTOPES OF LEAD AND MERCURY AND THE MASS DIFFERENCE D_2-He^4. Jay L. Benson,* Richard A. Damerow, and Richard R. Ries, School of Physics, University of Minnesota, Minneapolis, Minnesota (Received October 30, 1958).

Mass doublets have been measured from which atomic masses for the stable isotopes of mercury and lead are found. The doublet D_2-He^4 has also been measured and found to have a value 25.6074 ± 3 mMU.

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LIFETIMES OF THE FIRST EXCITED 0+ STATES OF Ca^{40} AND Zr^{90}. Robert M. Kloepper, Robert B. Day, and David A. Lind,* Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico (Received November 7, 1958).

The first excited states of Ca^{40} and Zr^{90} are 0+ states which decay to the 0+ ground state primarily by the emission of internal pairs or conversion electrons. These states have been excited by neutron inelastic scattering and their decay times have been observed by detecting the 0.51-Mev annihilation gamma rays as a function of time. The experimental method utilized a pulsed neutron source and the standard ring geometry often used for inelastic scattering measurements. Proton excitation of Ca^{40} in a different geometry was also used. To measure the lifetimes, the pulse-height spectra from a NaI(Tl) scintillation counter were recorded as a function of the delay time after the exciting pulse, and the intensity of the annihilation radiation was obtained from the area under the photopeak at 0.51 Mev. Selection of the time delays was made with the aid of either a time-to-pulse-height converter or a fast coincidence circuit. Mean lives for the first excited states of Ca^{40} and Zr^{90} were found to be $(3.4 \pm 0.2) \times 10^{-9}$ sec, respectively. Corresponding values of the reduced matrix elements ρ are 0.15 and 0.056.

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DIFFERENTIAL ELASTIC SCATTERING OF NEUTRONS FROM NEON. H. O. Cohn and J. L. Fowler, Oak Ridge National Laboratory, Oak Ridge, Tennessee (Received November 12, 1958).

Neutron elastic scattering from neon of natural abundance has been investigated in the energy range 0.8 to 1.9 Mev. The total cross section was measured by the standard method of attenuating a neutron beam by a sample, which in this case was neon gas contained in a meter long steel cylinder at high pressure. Differential cross sections were observed both by measuring neon recoil energies in a proportional counter and also by detecting neutrons with an anthracene crystal in coincidence with neon recoils. The combination of techniques shows up resonances at the following energies (in Mev) with spin and parity assignments given in parentheses: 0.91 (3/2 -); 1.28 (3/2 -); 1.31 (1/2 -); 1.37 (3/2 or 5/2 +); 1.62 (3/2 -); 1.68 (3/2 +); and 1.85 (...). These resonances have reduced widths of the order of one percent of the Wigner limits.

0-600 kev GAMMA-RAY SPECTRA FROM THERMAL NEUTRON CAPTURE IN THE REGION A =104 to 198. James E. Draper, Yale University, New Haven, Connecticut (Received October 28, 1958).

The energies and absolute intensities of prominent peaks in the 0-600 kev region of the gamma-ray spectrum following thermal neutron capture have been measured with a single NaI(Tl) scintillation spectrometer. The elements investigated were rhodium, silver, cadmium, indium, antimony, tellurium, iodine, samarium, europium, gadolinium, terbium, dysprosium, holmium, erbium, thulium, hafnium, tantalum, rhenium, iridium, platinum, and gold.

UNSTABLE PARTICLES AS TARGETS IN SCATTERING EXPERIMENTS. G. F. Chew, Radiation Laboratory, University of California, Berkeley,

California, and F. E. Low, Radiation Laboratory, University of California, Berkeley, California, and Department of Physics and Laboratory for Nuclear Science, Massachusetts Institute of Technology, Cambridge, Massachusetts (Received November 3, 1958).

A general method is suggested for analyzing the scattering of particle A by particle B, leading to three or more final particles, in order to obtain the cross section for the interaction of A with a particle which is virtually contained in B. Binding complications are absent if a plausible assumption about the location and residues of poles in the S-matrix is accepted. The method is useful for unstable particles from which free targets cannot be made; the special examples of pion and neutron targets are discussed in detail.

GEOMAGNETIC MEASUREMENTS ON HEAVY PRIMARY COSMIC RADIATION NEAR THE EQUATOR.

Robert E. Danielson,* School of Physics, University of Minnesota, Minneapolis, Minnesota (Received August 29, 1958).

The azimuth angle distribution of primary cosmic radiation with charges $Z \geq 6$ has been measured near the geomagnetic equator (at Guam) in horizontal emulsions with known orientation relative to the earth. The observed distribution is well described by using the centered dipole approximation to the earth's surface magnetic field (north pole at 79°N and 70°W), if the effect of the solid earth is included by assuming that the main cone is the allowed cone. This verifies the prediction that the penumbra becomes forbidden near the geomagnetic equator. The main cone equatorial cutoffs were applied to an assumed integral energy spectrum which is inversely proportional to the α power of the total energy per nucleon. The exponent α was found to be at least 1.82 ± 0.19.

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PHOTOPRODUCTION OF CHARGED π MESONS FROM NUCLEI.

John R. Waters,* Cornell University, Ithaca, New York (Received October 20, 1958).

The γ-ray beam from the Cornell synchrotron was used to investigate the photoproduction of low-energy charged π mesons from several nuclei at 35° to the beam direction. Maximum photon energies used were 800 Mev and 1 Bev. The yields of 40- and 80-Mev π^+ and 40-Mev π^- were observed with a detection system whose aperture was known absolutely. It was found that in some cases the meson yields were proportional to the atomic number, A, of the target rather than the more usual $A^{2/3}$. The data are compared with the predictions of an optical model of the nucleus. The deviations observed can be partially explained in terms of the inelastic scattering of high-energy mesons and of multiple meson production. An excitation function was measured for 80-Mev π^+ mesons from carbon by varying the machine energy; this indicated the presence of an appreciable number of multiply-produced mesons.

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EVIDENCE FOR THE TRANSITION OF A K^0 INTO A $\overline{K}^0$ MESON.

Frank S. Crawford, Jr., Marcello Cresti, Myron L. Good, Klaus Gottstein, Ernest M. Lyman, Frank T. Solmitz, M. Lynn Stevenson, and Harold Ticho, Radiation Laboratory, University of California, Berkeley, California (Received November 4, 1958).

Two pictures have been obtained in a liquid-hydrogen bubble chamber, each of which demonstrates the following sequence:

(1) $\pi^- + p \rightarrow Y^0 + K^0$, $[Y^0_A = \Lambda^0,\ Y^0_B = \Sigma^0]$

(2) $\Lambda \rightarrow p + \pi^-$,

(3) $K^0 \rightarrow (\sim 50\%) K^0 + (\sim 50\%) \overline{K}^0$,

(4) $\overline{K}^0 + p \rightarrow \Sigma^+ + \pi^0$,

(5) $\Sigma^+ \rightarrow n + \pi^+$,

where step (3) is not directly observable, but is a prediction of the theory of Gell-Mann and Pais. On the basis of two events, the cross section for process (4) is 50 mb.

FLUX AND ENERGY SPECTRUM OF COSMIC-RAY α-PARTICLES DURING SOLAR MAXIMUM.

P. S. Freier, E. P. Ney, and C. J. Waddington, University of Minnesota, Minneapolis, Minnesota (Received November 10, 1958).

The fluxes of primary cosmic-ray α-particles over Minnesota and Texas have been measured during the present period of maximum solar ac-

tivity. A value of 136 ± 9 α-particles/m^2 sec sterad was measured over Minnesota and of 68 ± 4 α-particles/m^2 sec sterad over Texas. In both cases these values are significantly lower than those observed at solar minimum. The energy spectrum of these particles has been determined between 200 Mev/nucleon and 3.0 Bev/nucleon. It is shown that the slope of the integral spectrum is less than that observed at solar minimum and that a significant number of low-energy particles is still present. A possible mechanism for these changes is discussed briefly. The determination of energies of particles from a measurement of their ionization is discussed in detail in an Appendix.

LIFETIMES AND DECAY SPECTRA OF τ'^+ AND $K_{\mu 3}$. S. Taylor,* G. Harris,† J. Orear,‡ J. Lee, and P. Baumel, Columbia University, New York, New York (Received November 13, 1958).

Nuclear emulsion stacks were exposed to the Bevatron ~360-Mev/c unseparated K^+ beam at two different distances from the target. The K_L to τ' and K_L to $K_{\mu 3}$ ratios at the two positions combined with the known K_L lifetime yield $(1.0^{+0.5}_{-0.3}) \times 10^{-8}$ sec for the τ' mean life and $(1.2^{+1.0}_{-0.4}) \times 10^{-8}$ sec for the $K_{\mu 3}$ mean life. These lifetimes, when compared with other K^+ lifetimes, are consistent with a single lifetime for all K^+ mesons. The π^+ energy spectrum from τ'^+ decay is obtained, using 72 events, and by Dalitz-type analysis is found to be consistent with spin zero but would be unlikely for spin one. Hence, the τ' spin is consistent with being the same as that of the τ. The μ^+ energy spectrum from $K_{\mu 3}{}^+$ decay is obtained from 0 to 63 Mev, using 47 events, and is found to be consistent with the energy distribution of the density of final states in the region. The relative abundances of τ^+, τ'^+, and $K_{\mu 3}{}^+$ are found to be 0.052 ± 0.003, 0.015 ± 0.002, and 0.028 ± 0.004, respectively, the $K_{\mu 3}$ abundance having been estimated by assuming that the μ^+ spectrum above 63 Mev is proportional to the density of final states.

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EFFECT OF RECOIL ON THE ELASTIC SCATTERING OF HIGH-ENERGY ELECTRONS BY ZERO-SPIN NUCLEI. Leslie L. Foldy, Case Institute of Technology, Cleveland, Ohio, and Kenneth W. Ford,* Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico, and Donald R. Yennie, University of Minnesota, Minneapolis, Minnesota (Received October 29, 1958).

The effect of nuclear recoil on the elastic scattering of high-energy electrons or muons by zero-spin nuclei is studied by adapting the Breit two-particle Hamiltonian to the case that one of the two particles is of finite size, is spinless, and is nonrelativistic, the other being a normal point Dirac particle. A radial and angular separation of the Dirac equation is still possible. To leading order in the parameter (electron energy)/(nuclear mass), the effect of the dynamic recoil terms is to rotate the scattering amplitude vectors in the complex plane without changing their magnitudes, a result which is independent of the shape and size of the nuclear charge distribution. To this order, the cross section is affected only by the kinematic recoil corrections. The dynamic recoil terms also influence the scattering amplitudes through terms of order (electron mass)/(nuclear mass). These corrections, owing to large amplification factors in going from phase shifts to cross section, may be of some significance in muon scattering, but are probably of no importance in the analysis of high-energy electron scattering. The dynamic effect is proportional to nuclear charge and therefore nearly as great for heavy as for light nuclei.

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SOLUBLE PROBLEM IN DISPERSION THEORY. M. L. Goldberger and S. B. Treiman, Palmer Physical Laboratory, Princeton University, Princeton, New Jersey (Received November 12, 1958).

The Lee model is modified by addition of a new field θ' and a weak coupling $N+\theta \rightarrow N+\theta'$, which leads to instability of the V particle: $V \rightarrow N+\theta \rightarrow N+\theta'$. The decay amplitude is calculated to lowest order in the weak coupling by dispersion relation methods. In effect we are required to study a set of simultaneous dispersion relations. The problem is completely soluble and serves to

clarify the essential structure of dispersion methods. The results agree with what one obtains, more easily in the present case, by direct methods.

SPECTRAL REPRESENTATIONS IN PERTURBATION THEORY. II. TWO-PARTICLE SCATTERING. Robert Karplus, Charles M. Sommerfield, and Eyvind H. Wichmann, Radiation Laboratory and Physics Department, University of California, Berkeley, California (Received November 7, 1958).

In this paper a particular term in the perturbation expansion for the two-particle scattering amplitude is examined. We consider the real plane defined by the square of the total four-momentum and the square of the momentum transfer, and show that the scattering amplitude is an analytic function of both variables in certain connected region in this plane. The precise boundary of the region is found. The purpose of this work is to find some conditions that integral representations of the scattering amplitude must satisfy, with the hope that such examples may aid the study of such integral representations in general.

We also apply our general result to some particular cases of physical interest.

FERMI INTERACTION CAUSED BY INTERMEDIARY CHIRAL BOSON. Yasutaka Tanikawa,* Institute for Advanced Study, Princeton, New Jersey, and Satosi Watanabe, International Business Machines Research Laboratory, Ossining, New York (Received May 22, 1958).

The Fermi interaction is interpreted as a nonlocal interaction resulting from a double Yukawa-type interaction in which the intervening boson has a definite chirality. The theory is quantizable and renormalizable, and in the "local" limit, the results agree with the usual (V-A) theory of the direct Fermi interaction. The general framework of the theory not only gives a basis for the existence of parity-nonconserving interactions, but also determines the allowed forms of such interactions. The nonlocal effect of the intervening boson propagator tends to give an improved agreement with experiments.

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CONSEQUENCES OF RENORMALIZABLE WEAK INTERACTIONS. S. Oneda* and Y. Tanikawa,† Institute for Advanced Study, Princeton, New Jersey (Received June 2, 1958).

Consequences of the theory of renormalizable weak interactions caused by the intermediary chiral bosons are further investigated with regard to the decay processes of the pion, K-meson, and hyperons. It is shown that all those aspects of weak couplings which can be satisfactorily explained in terms of the local Fermi interaction with V-A combination can be equally well reproduced by the present model. Furthermore, on account of the convergence of the present theory, the results of computation are less ambiguous than the local Fermi interaction model. Finally, some possible experimental tests of the theory, mainly by the use of the nonlocalizability of the lepton interactions innerent in this theory, are discussed.

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